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Classical and Quantum Aspects of the Color Glass Condensate

March 7 – 11, 2005

Organizers:

Dmitri Kharzeev, Anna Stasto, Kirill Tuchin, Werner Vogelsang

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Preface to the Series

The RIKEN BNL Research Center (RBRC) was established in April 1997 at Brookhaven National Laboratory. It is funded by the "Rikagaku Kenkyusho" (RIKEN, The Institute of Physical and Chemical Research) of Japan. The Center is dedicated to the study of strong interactions, including spin physics, lattice QCD, and RHIC physics through the nurturing of a new generation of young physicists.

The RBRC has both a theory and experimental component. At present the theoretical group has **4** Fellows and **3** Research Associates as well as **11** RHIC Physics/University Fellows (academic year **2003-2004).** To date there are approximately 30 graduates from the program of which 13 have attained tenure positions at major institutions worldwide. The experimental group is smaller and has **2** Fellows and **3** RHIC Physics/University Fellows and 3 Research Associates, and historically 6 individuals have attained permanent positions.

Beginning in **2001** a new **RIKEN** Spin Program (RSP) category was implemented at RBRC. These appointments are joint positions of RBRC and RIKEN and include the following positions in theory and experiment: RSP Researchers, RSP Research Associates, and Young Researchers, who are mentored by senior RBRC Scientists. A number of RIKEN Jr. Research Associates and Visiting Scientists also contribute to the physics program at the Center.

RBRC has an active workshop program on strong interaction physics with each workshop focused on a specific physics problem. Each workshop speaker is encouraged to select a few of the most important transparencies from his or her presentation, accompanied by a page of explanation. This material is collected at the end of the workshop by the organizer to form proceedings, which can therefore be available within a short time. To date there are sixty-eight proceeding volumes available.

The construction of a 0.6 teraflops parallel processor, dedicated to lattice QCD, begun at the Center on February 19,1998, was completed on August **28,** 1998 and is still operational. A 10 teraflops QCDOC computer in under construction and expected to be completed this year.

N. P. Samios, Director
November **2004**

*Work performed under the auspices of **U.S.D.O.E.** Contract **No. DE-AC02-98CH10886.**

CONTENTS

Introduction and overview

The high energy limit of Quantum Chromodynamics is one of the most fascinating areas in the theory of strong interactions. Over a decade ago the HERA experiment at DESY in Hamburg provided strong evidence for the rise of the proton structure function at small values of the Bjorken variable x. This behavior can be explained as an increase of the gluon density of the proton with energy or correspondingly with smaller values of x. This increase can be attributed on the other hand to the large probability of gluon splitting in QCD. The natural fiamework for describing the gluon dynamics at small x is the Balitskii-Fadin-Kuraev-Lipatov formalism developed some 30 years ago. It predicts that the gluon density grows very fast with increasing energy, as a power with a large intercept. This increase has to be tamed in order to satisfy the unitarity bound. Over two decades ago, Gribov, Levin and Ryskin proposed the mechanism called the parton saturation, which slows down the fast rise of the gluon density. This formalism accounts for an additional gluon recombination apart from the pure gluon splitting. It leads to the very interesting non-linear modification of the evolution equations for the gluon distributions. Since then, much progress has been made in the theoretical formulation of the parton saturation. Currently the most complete theory for parton saturation is the Color Glass Condensate (CGC) with the corresponding renormalization group functional evolution equation, the JIMWLK equation, which describes the nonlinear evolution of the gluon density at small values of x and in the regime where gluon fields are strong. The simpler form of the JIMWLK equation, the Balitskii-Kovchegov (BK) equation has been successfully used to explain the experimental data on proton structure function. The models, which include the parton saturation, have been applied to explain the experimental data at Tevatron and RHIC. In the latter case the Color Glass Condensate can be thought of as an initial stage for the subsequent formation of the Quark Gluon Plasma. Despite its success in describing various observables, the parton saturation phenomenon still needs deeper understanding and improvements; in particular, the existence or limitations on geometrical scaling, the edge effects in the high energy collisions, or impact parameter dependence. In particular it has been recently realized that the current evolution equations of CGC, the JIMWLK equations miss some of the important contributions coming from the resummation of the so-called Pomeron loops. These terms are known to provide sizeable corrections to the asymptotic high energy behavior. Also, the CGC formalism was constructed within the leading logarithmic approximation, and it is known that the corrections which go beyond this order are large.

This aim of this workshop was to bring together experts in the field to study and discuss the Color Glass Condensate theory and related topics. Larry McLerran gave an introduction to the theory of Color Glass Condensate and reviewed the current status of the parton saturation. He discussed also the relation between the CGC and the Quark Gluon Plasma, in particular the problems related with rapid thermalization and the flow in ultrarelativistic heavy ion collisions. The Pomeron effective theory and the fluctuation–saturation duality was discussed by Kazunori Itakura. He presented the new evolution equations of Color Glass Condensate theory that take into account the Pomeron loops. The improved Hamiltonian of the new evolution equation, which aims to describe the weak and strong fields regime, is likely to possess the intriguing self-duality property.

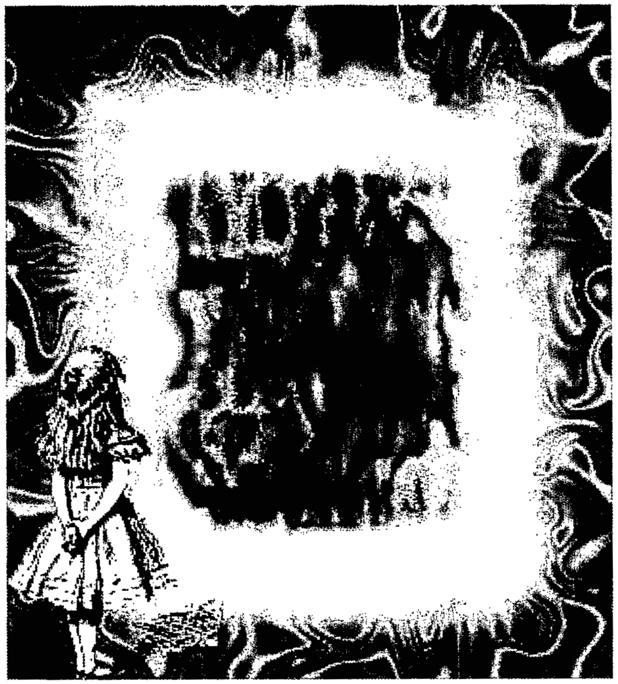
Boris Kopeliovich discussed the phenomenological aspects, namely the application of the saturation models to ultrarelativistic heavy ion collisions. He carefully studied the kinematical conditions in which the Color Glass Condensate description is expected to be valid and pointed out other dynamical effects that can occur in the same regime. An interesting relation between the QCD at high energies and the statistical physics was presented by Alfred Mueller. The large gluon occupation numbers suggest that one can use tools of statistical physics to describe the evolution at high energies in QCD. In particular the role of fluctuations at the beginning of the evolution, namely in the dilute regime, is expected to be very large. Therefore, the reformulation of the JIMWLK equation in order to properly account for these effects is needed. Heribert Weigert gave a nice overview of the current status of the solutions to the BK and JIMWLK equations. He showed that the running coupling corrections play an important role and need to be included in the evolution. He also discussed an intriguing observation of the formal similarity of the evolution equation for the non-global jet observables and the BK equation at small x. The investigation of the Odderon evolution equation within the framework of the Color Glass Condensate was presented by Yoshitaka Hatta. He presented the derivation of the evolution equations for the amplitudes describing the odderon exchanges between the Color Glass Condensate and the two types of the projectiles: a color dipole and a system consisting of three quarks. He showed that in the linear regime the equations reduce to the Bartels-Kwiecinski-Praszalowicz evolution equation. Stephen Wong discussed recent developments that lead to the generalized JIMWLK evolution equation including the Pomeron loops. It has been recently realized that the JIMWLK equation has certain deficiencies, namely, that it only includes the non-linear effects due to the Pomeron mergings but misses important contributions coming from the Pomeron splittings. An extended version of the JIMWLK equation has been formulated which cures this problem and it includes the Pomeron loops. The current status of the reggeon field theory has been presented by Jochen Bartels. He discussed in detail the existing ingredients of this theory, which are the reggeized gluon and the vertex functions. He discussed the relation and differences with respect to the Color Glass Condensate theory. Alex Kovner discussed the extension of the generalization to JIMWLK presented by Wong by including yet higher order functional derivatives in the JIMWLK equation. This leads to the emergence of the self-dual theory, in which the projectile and the target are treated symmetrically. Francois Gelis presented the calculation of the quark-antiquark production cross section in the CGC formalism in the proton-nucleus collisions. He showed that generally the high energy factorization is violated by the presence of the saturation scale in the problem. Carlos Salgado presented numerical solutions to the Balitsky-Kovchegov equation and discussed their applications to the description of the deep inelastic scattering collisions of lepton-proton and lepton-nucleus. He showed that the saturation model with geometrical scaling leads to the very good description of various observables in deep inelastic and nucleus-nucleus collisions. Ian Balitsky developed an approach in which the high energy scattering in QCD can be viewed as a scattering of two shock waves. He presented the Wilson-line functional integral for effective action that contains all the information about the high energy scattering in the leading logarithmic .approximation. Jianwei Qiu discussed the transition from the parton model to parton saturation. He showed that in the case of the standard DGLAP evolution, the resummation of the dynamical power corrections leads to the shift of the parton momentum fraction by a single parameter. Elena Ferreiro described a phenomenological model of color strings for the soft dynamics in QCD. In this model the color strings are small areas in transverse space filled with color field created by the colliding partons. Michael Lublinsky presented a probabilistic approach to the description of the high energy QCD evolution. He showed a functional evolution equation that, accommodates the nonlinear dynamics. The problem of the thermalization in heavy ion collisions was discussed by Yuri Kovchegov. He demonstrated that at any order of the perturbative expansion the gluon field generated in the ultrarelativistic heavy ion collision leads to the energy density, which scales as an inverse proper time. This has to be contrasted with the hydrodynamics-driven expansion of the quark-gluon plasma which leads to the energy density, which scales as a higher power of the inverse proper time. Ismail Zahed discussed the RHIC fireball production in a theoretical framework of the AdS/CFT correspondence. Adrian Dumitru talked about the observational constraints on the saturation scale from cosmic ray airshower data. The simulations at highest energies of cosmic rays indicate that there is a substantial sensitivity to the QCD evolution scenario. There are indications that the saturation scale grows at a slower rate than predicted by HERA or RHIC data.

Peeking through the Colored looking Glass

Larry McLerran
Physics Department
Brookhaven National Laboratory
Upton, NY 11973

Peeking through the Colored Looking Glass

A perspective on Future Directions



Color Glass Condensate as a Media

Whatever-ons:
Little wiggles of the CGC
Pomerons, Odderons, Reggeons

Ploops: (Pomeron loops)
How a little fluctuation becomes a big problem

The CGC and the QGP:
Is the sQGP really the CGC?
Is rapid "thermalization" due to the CGC?
Does flow arise largely from the CGC?

Comments about the LHC:
The CGC Machine

Reggeons, Pomerons and Odderons

Reggeons:

Mathematical objects which which allow the computation of scattering of hadrons. Found in complex angular momentum analysis of scattering matrix

Pomeron:

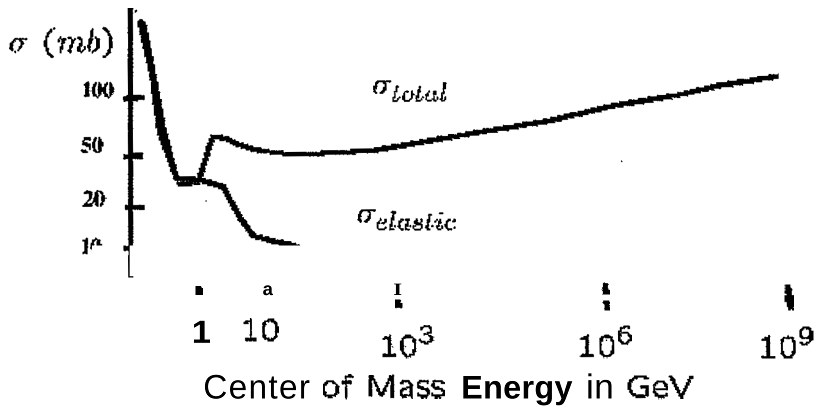
That Reggeon which controls the total cross section at high energy.
Universal dependence of energy at high energy.
Imaginary part of T matrix

Odderon:

Pomerons peculiar brother
Real part of T matrix at high energy

The Pomeron: A Modern Perspective

The total hadronic cross section:

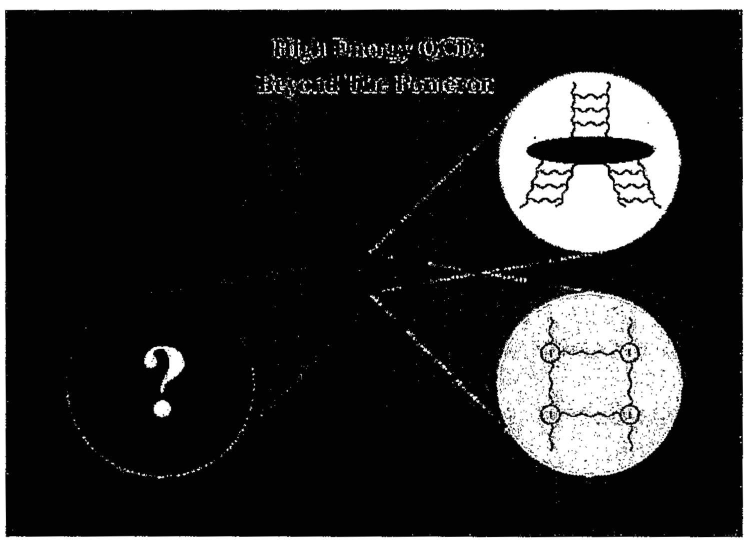


Describes total cross section
Original version

$$\sigma_{tot} - constant$$

After growing cross sections:

$$|A|_{bare} \sim E^{\delta}$$



Pomeron: Vacuum quantum numbers

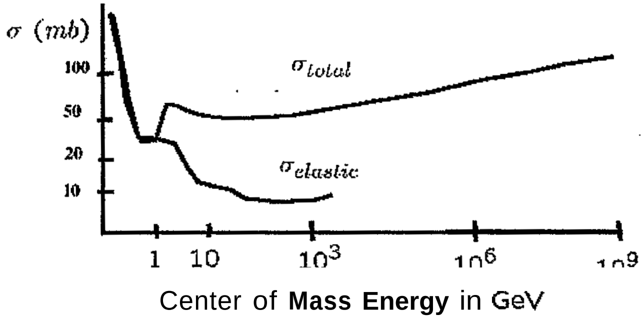
Two gluon exchange

Bare Pomeron related to growth of gluon densities

Saturation <=> High density of pomerons

The Growth of Gluon Density Explains Slow Growth of Total Cross Section

The total hadronic cross section:



Transverse distribution of gluons:

$$\frac{dN}{dyd^2r_T} = Q_{sat}^2(y)e^{-2m_\pi r_T}$$

Transverse profile set by initial conditions

Size is determined when probe sees a fixed number of particles at some transverse distance

$$e^{\kappa y}e^{-2m_\pi r_T} \sim constant$$

$$\sigma \sim r_T^2 \sim y^2 \sim ln^2(E/\Lambda_{QCD})$$

The Pomeron and the CGC

$$\int_{\Lambda^+} [dA][d\rho] e^{iS[A,\rho]-F[\rho]}$$

Separation between fast and slow degrees of freedom => Renormalization group

$$Z = e_{-}^{F[\rho]}$$

Weight for source fluctuations

$$\frac{d}{d\eta} Z = -H(\rho, d/d\rho) Z$$

$$\eta = ln(1/x)$$

JIMWLK Equation:

H is Hamiltonian with

No potential, assumes strong fields

For weak fields:

$$H = \int [dx][dy][dz] \frac{(x-y)^2}{(x-z)^2(y-z)^2} \frac{d}{d\rho(x)} \{\rho(x) - \rho(z)\}\{\rho(y) - \rho(z)\} \frac{d}{d\rho(y)}$$

Pomeron: Weak field excitation $\sim \rho(x) \; {}_{+} \; \rho(y)$

Pomeron: Saturation effects

$$O(x,y) = < tr(U(x)U^\dagger(y)) >$$

$$\frac{d}{d\eta}O(x,y) = \kappa\alpha_s \int [dz] \frac{(x-y)^2}{(x-z)^2(y-z)^2} \{O(x,z) + O(z,y) - O(x,y) - O(x,z)O(z,y)\}$$

The real part of O is the Pomeron amplitude
Balitsky-Kovchegov equation

BK equation has exponential growth in y for transverse momentum scales greater than the saturation momenta; Power law growth for momenta less than the saturation moments.

The saturation momentum never saturates

$$Q_{sat}(y) \sim e^{by}$$

Fluctuation-saturation duality and Pomeron effective theory

K. Itakura

Service de Physique Théorique, CEA/Saclay, F-91191 Gif-sur-Yvette, France

Abstract

We propose an effective theory which governs Pomeron dynamics in QCD at high energy, in the leading logarithmic approximation, and in the limit where N_c, the number of colors, is large. In spite of its remarkably simple structure, this effective theory generates precisely the evolution equations for scattering amplitudes that have been recently deduced from a more complete microscopic analysis. It accounts for the BFKL evolution of the Pomerons together with their interactions: dissociation (one Pomeron splitting into two) and recombination (two Pomerons merging into one). It is constructed by exploiting a duality principle relating the evolutions in the target and the projectile, more precisely, splitting and merging processes, or fluctuations in the dilute regime and saturation effects in the dense regime. The simplest Pomeron loop calculated with the effective theory is free of both ultraviolet or infrared singularities.

This talk is based on the paper hep-ph/0502221 by J.P.Blaizot, E,Iancu, K.Itakura, and D.Triantafyllopoulos.

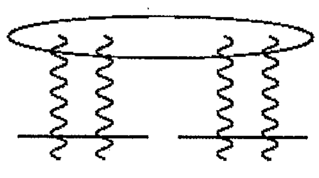
Introduction (2/3)

More precisely,

- deep understanding about the difference between BK and Balitsky eq.
- analogy of gluon dynamics (multiplication vs recombination) with statistical physics (reaction-diffusion dynaimcs) described by FKPP eq.

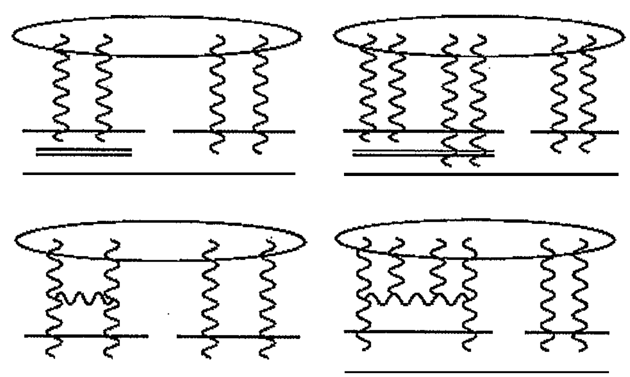
→ BK-JIMWLK hierarchy fails to correctly describe fluctuation phenomena.
[Iancu-Triantafyllopoulos]

→ This becomes manifest when we consider dipole-pair scattering



Tree dipole pair diagram

(Described by CDP)



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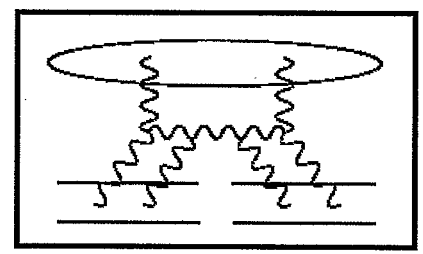
Described by JIMWLK



Not included in JIMWLK

Duality (2/4)

Dipole-dipole scattering in the dilute regime

Scattering amplitude in the eikonal approximation

$$\langle T(\boldsymbol{x},\boldsymbol{y})\rangle_Y = 1 - \frac{1}{N_c}\left\langle \mathrm{tr}(V_{\boldsymbol{x}}^\dagger V_{\boldsymbol{y}})\right\rangle_Y \qquad V_{\boldsymbol{x}}^\dagger[\alpha] = \mathrm{P}\exp\left(ig\int dx^- \alpha^a(x^-,\boldsymbol{x})t^a\right)$$

$$\simeq T_Y^0(\boldsymbol{x},\boldsymbol{y}) \equiv \frac{g^2}{4N_c}\left\langle\left(\alpha^a(\boldsymbol{x})-\alpha^a(\boldsymbol{y})\right)^2\right\rangle_Y$$ weak-field limit (single scatt)

can be equivalently rewritten as (to the order of interest, 2nd order)

$$T_Y^0(\boldsymbol{x},\boldsymbol{y}) = \int \mathrm{D}[\alpha_R]\; W_Y[\alpha_R]\left\langle 1 - \mathrm{e}^{\,i\int d^2\boldsymbol{z}\,\rho_L^a(\boldsymbol{z})\alpha_R^a(\boldsymbol{z})}\right\rangle_Q$$

average over target average over projectile

with $\rho_L^a(\boldsymbol{z}) \equiv Q_L^a\left[\delta^{(2)}(\boldsymbol{z}-\boldsymbol{x}) - \delta^{(2)}(\boldsymbol{z}-\boldsymbol{y})\right]$ $\quad \langle Q_L^a\rangle_Q = 0, \quad \langle Q_L^a Q_L^b\rangle_Q = \delta^{ab}\frac{g^2}{2N_c}$

Scattering matrix (projectile and target have rapidities *y* and *Y-y*)

$$\boxed{S_Y = \int \mathrm{D}[\alpha_R]\; W_{Y-y}[\alpha_R]\int \mathrm{D}[\alpha_L]\; W_y[\alpha_L]\;\mathrm{e}^{\,i\int d^2\boldsymbol{z}\,\rho_L^a(\boldsymbol{z})\alpha_R^a(\boldsymbol{z})}}$$

Iancu-Mueller factorization formula NPA730(2004)

Duality (3/4)

Lorentz invariance of the S-matrix → S_Y must be independent of y

[Kovner-Lublinsky]

$$0 - \frac{\partial S_Y}{\partial y} \quad \int D[\alpha_R] \int D[\alpha_L]\; e^{i\int d^2z\, \rho_L^a(z)\alpha_R^a(z)}$$

$$\left\{ \left(\frac{\partial}{\partial y} W_{Y-y}[\alpha_R]\right) W_y[\alpha_L] + W_{Y-y}[\alpha_R] \left(\frac{\partial}{\partial y} W_y[\alpha_L]\right) \right\}$$

Backward evolution
In the target

$$-H\left[\alpha_R, \frac{\delta}{i\delta\alpha_R}\right] W_{Y-y}[\alpha_R]$$

Normal evolution
In the projectile

$$\cdot H\left[\alpha_L, \frac{\delta}{i\delta\alpha_L}\right] W_y[\alpha_L]$$

$$H\left[\alpha_R, \frac{\delta}{i\delta\alpha_R}\right] e^{i\int d^2z\, \rho_L^a(z)\alpha_R^a(z)} - H\left[\frac{\delta}{i\delta\rho_L}, \rho_L\right] e^{i\int d^2z\, \rho_L^a(z)\alpha_R^a(z)}$$

Self-duality condition

$$\boxed{H\left[\alpha_L, \frac{\delta}{i\delta\alpha_L}\right] W_y[\alpha_L] - H^\dagger\left[\frac{\delta}{i\delta\rho_L}, \rho_L\right] W_y[\rho_L]}$$

Pomeron effective theory (1/3)

One can use duality to construct an effective theory of interacting Pomerons (perturbative, bare) in the dilute regime:

"Pomeron" = $T_0(\boldsymbol{x},\boldsymbol{y}) \equiv \dfrac{g^2}{4N_c}[\alpha^a(\boldsymbol{x}) - \alpha^a(\boldsymbol{y})]^2$

"κ Pomerons" = $T_0^{(\kappa)}(\boldsymbol{x}_1,\boldsymbol{y}_1,...,\boldsymbol{x}_\kappa,\boldsymbol{y}_\kappa) = T_0(\boldsymbol{x}_1,\boldsymbol{y}_1)...T_0(\boldsymbol{x}_\kappa,\boldsymbol{y}_\kappa)$

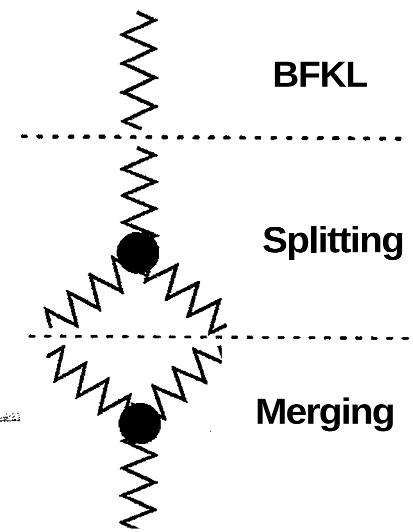
The total Hamiltonian is self-dual:

$$\boxed{H^\dagger = H_0^\dagger + H_{1\to 2}^\dagger +}$$

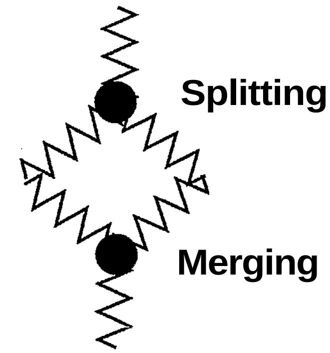
$H_0^\dagger$ = "BFKL" from weak-field exp. of JIMWLK, "free" part without number changing int, self-dual by itself

$H_{1\to 2}^\dagger$ = "Splitting": important in the dilute regime [Mueller,Shoshi,Wong]

$H_{2\to 1}^\dagger$ = "Merging": dual of splitting



Application



Minimal Pomeron loop from 2 step evolutions of one Pomeron $\mathbb{PL} = H^\dagger_{1\to 2} H^\dagger_{2\to 1} T_0$

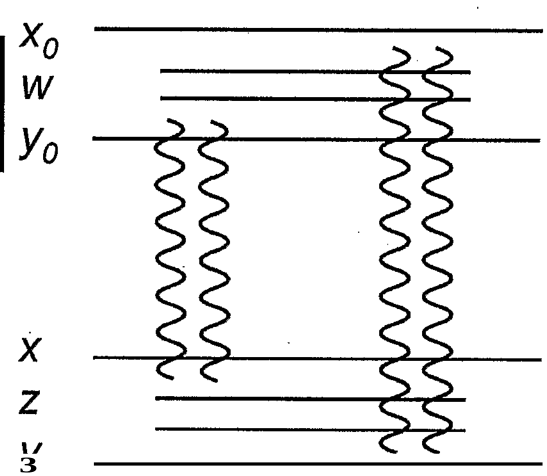
For *dipole-dipole* scattering (x,y) and (x_0,y_0)

$$\mathbb{PL}^0 = -2\left(\frac{\bar{\alpha}_s}{2\pi}\right)^2 \alpha_s^4 \int_{z,w} \mathcal{M}(x,y,z)\mathcal{M}(x_0,y_0,w)\mathcal{A}_0(x,z|x_0,w)\mathcal{A}_0(z,y|w,y_0).$$

Minus sign
→contribute to saturation

Two evolutions
two dipole kernels

Two gluon exchange amplitude of dipole-dipole scattering is $\alpha_s^2 \mathcal{A}_0$



z and w integrations are IR and UV finite

Effects which can mimic the CGC

Boris Kopeliovich

CGC Workshop
March 7, 2005

Conditions to be watched searching for a signal of CGC

- x_2 must be sufficiently small to provide a longitudinal overlap of gluons originated from different nucleons in a row.

  There are effects at large x_2 which might be misinterpreted as CGC

- This is not an easy-to-fulfil condition. The gluonic fluctuations are heavy and shortliving.

- Another difficult-to-fulfil condition is an overlap in impact parameters. Gluons are located within small spots with area an order of magnitude smaller than the proton.

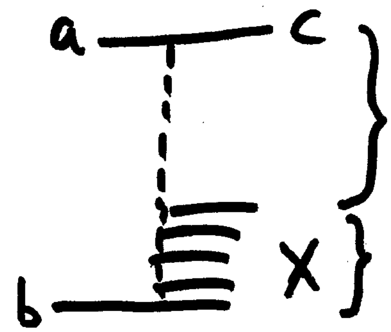
- One should watch x_1: even if x_2 is small, coherence vanishes at large x_1.

$$x_2^{eff} \gtrsim \frac{x_2}{1-x_1}$$

- Lessons from J/ψ production and other reactions at large x_1.
- Cronin effect in the proton fragmentation region. BRAHMS data.

Any reaction, $a+b \to c+X$, ($c = h, e\bar{e}, J/\psi, \ldots$) is a large rapidity gap (LRG) process at $x_F \longrightarrow 1$

Rapidity intervals



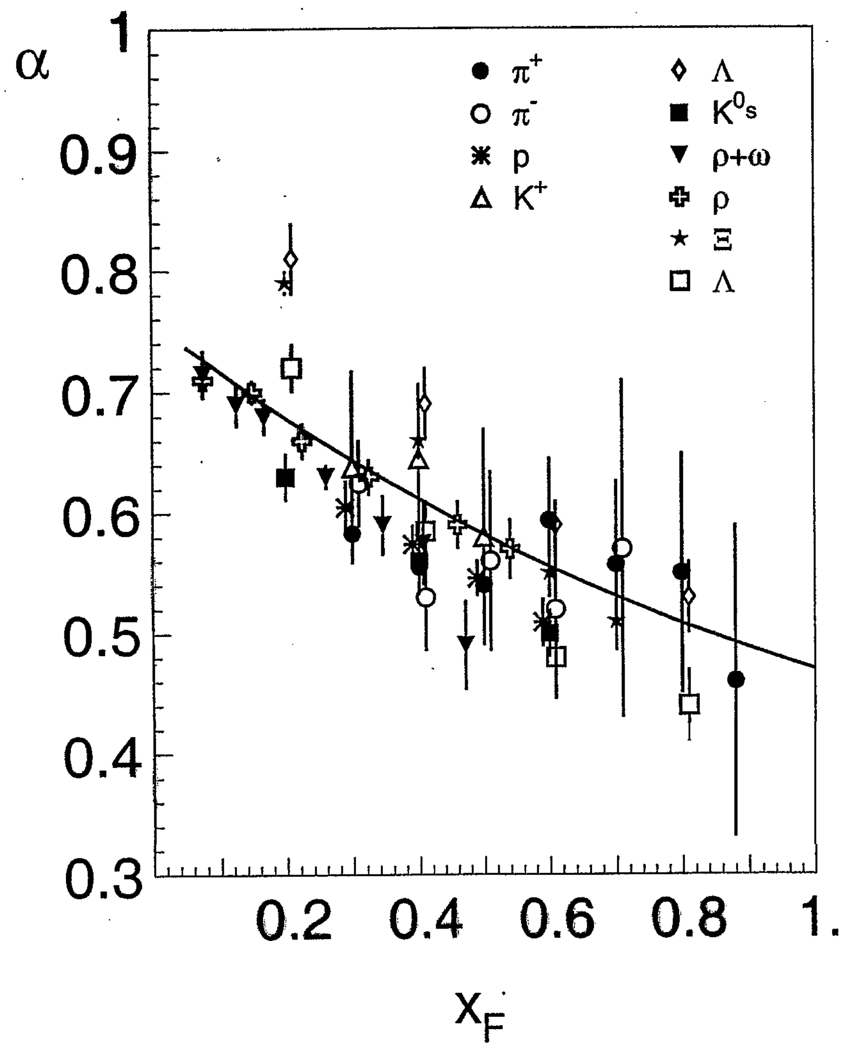
$\ln \frac{s}{M^2} = \ln \frac{1}{1-x_F}$

$\ln \frac{M^2}{s_0}$

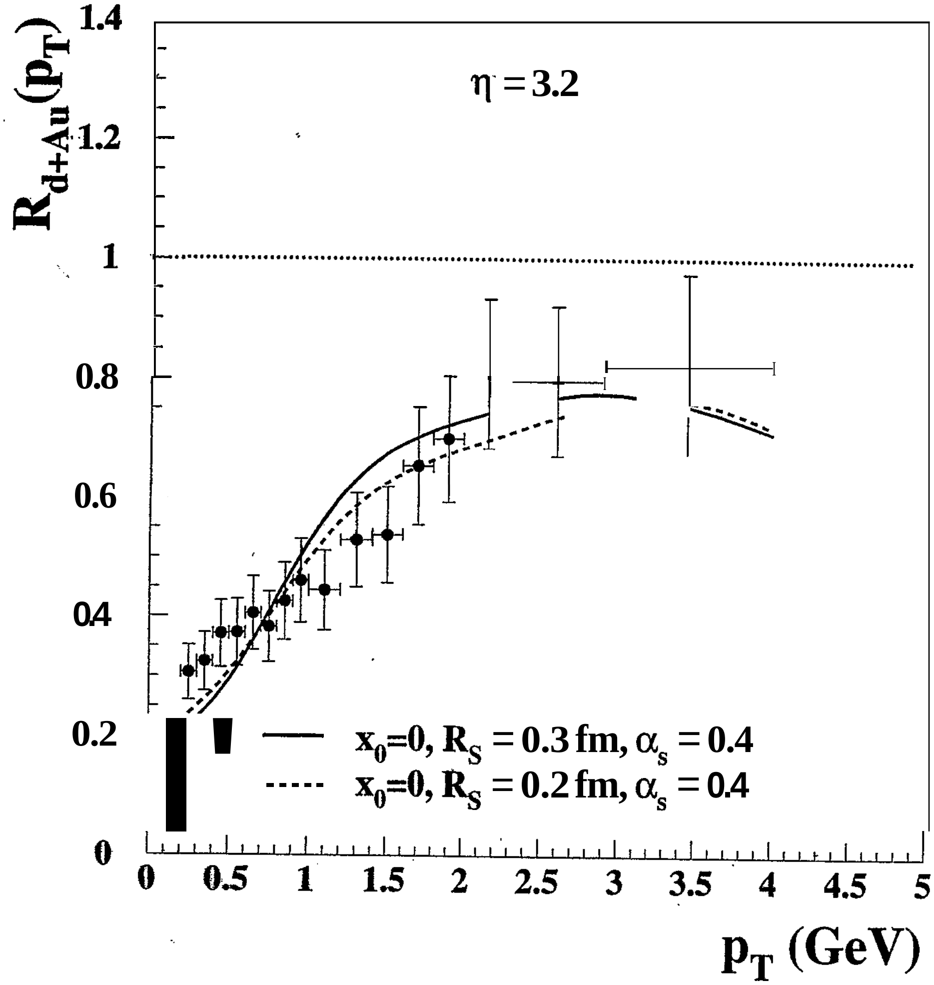
The probability to radiate no gluons in the rapidity interval $\Delta y = \ln \frac{1}{1-x_1}$ is suppressed by the Sudakov's formfactor $S(\Delta y)$ which violates QCD factorization

$$\sigma(pA \to hX) \propto A^{\alpha}$$

Different hadrons at different energies ($70 < E < 400$ GeV) are suppressed same way.



J. Nemchik
I. Potashnikova
M. Johnson
I. Schmidt
B.K.



Of Colored Glass & Jets in Medium

QCD @ high energies in heavy ion collisions

Heribert Weigert
Universität Regensburg

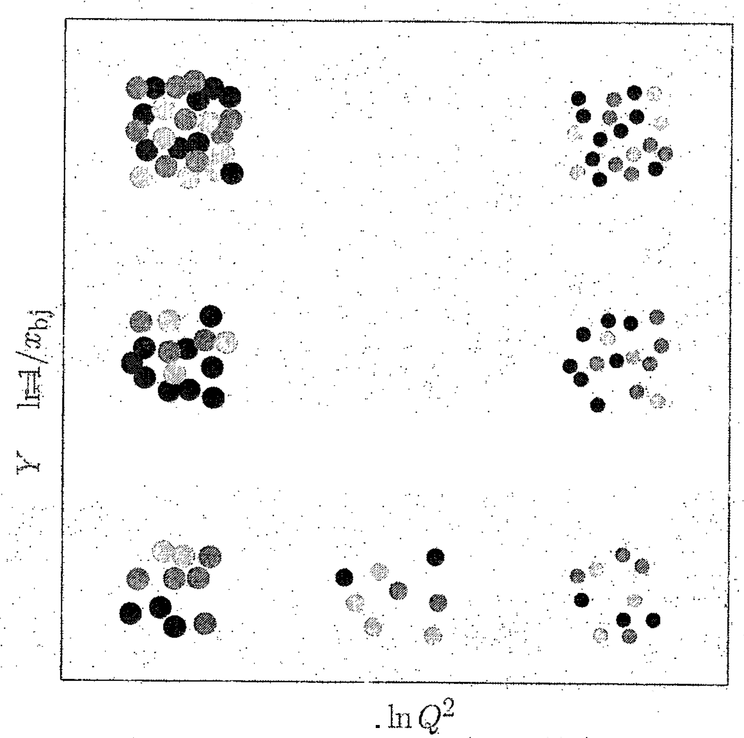
2005

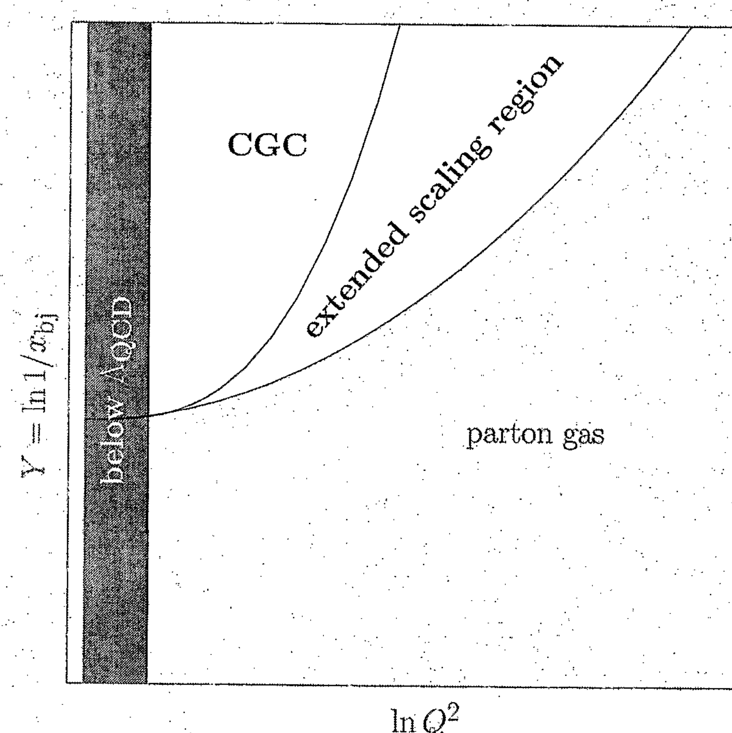
Towards pQCD beyond leading twist inclusive settings

Take away leading twist

- Example: Color Glass Condensate (CGC) @ small x

towards small x at fixed Q^2





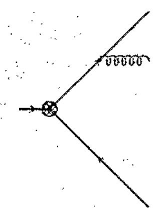
- Example: Jets in a Medium (LHC); non-global jet observables

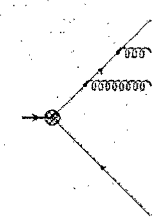
Towards pQCD beyond leading twist inclusive settings

Theoretical tools:

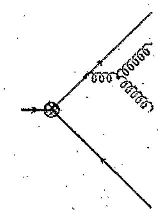
Common feature: soft gluon radiation @ high energies

- photon like contributions





- QCD: charged gluons



- enhanced by phase space integrals $\frac{dE}{E}\frac{d\theta}{\theta} \Longrightarrow \alpha_s \ln E \ln \theta$
- all orders calculation needed $\sum_{n=0}^{\infty} (\alpha_s \ln E)^n \ldots$
- gluons charged $\Longrightarrow$ radiation nonlinear in QCD

- Evolution equations:
  JIMWLK (CGC) Jet analogues thereof

The JIMWLK evolution equation

explicit form Heribert Weigert **Nucl. Phys. A703**, 2002, '823

$$\frac{d}{d\eta} Z_\eta[U] = -H_{\mathsf{JIMWLK}}[U] \quad Z_\eta[U]$$

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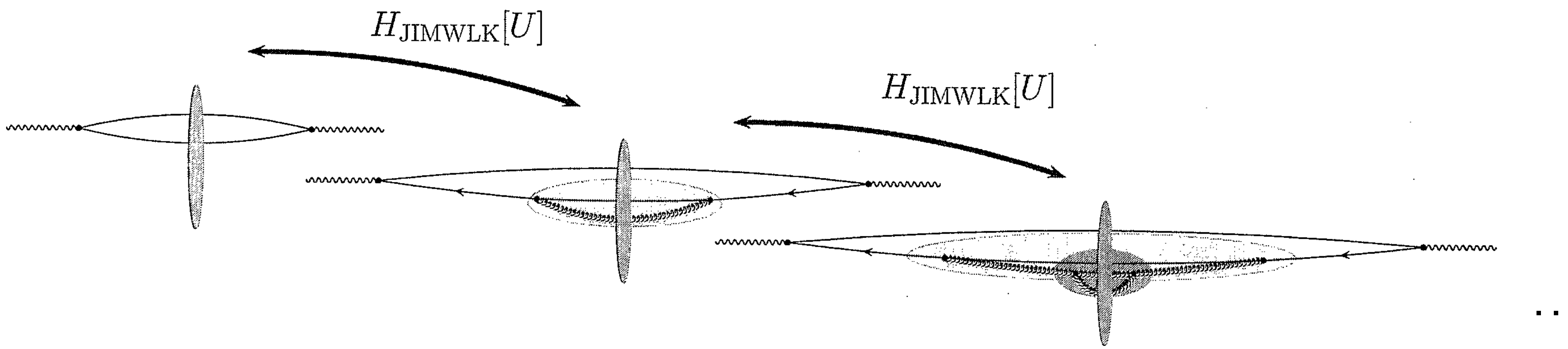
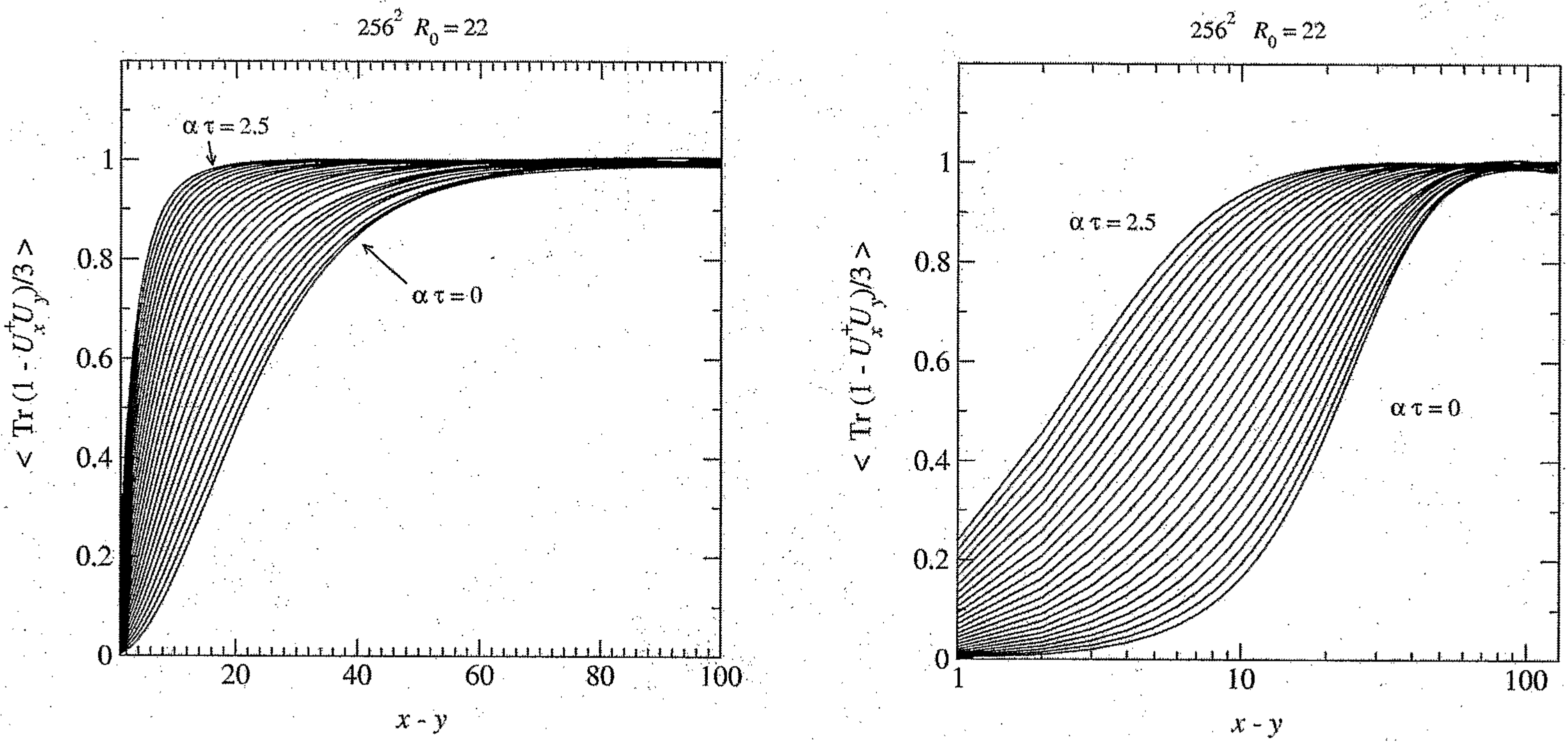


➡ energy dependence of $\langle \ldots \rangle(\eta)$

JIMWLK: simulations

Simulations show scaling (Rummukainen & H.W)

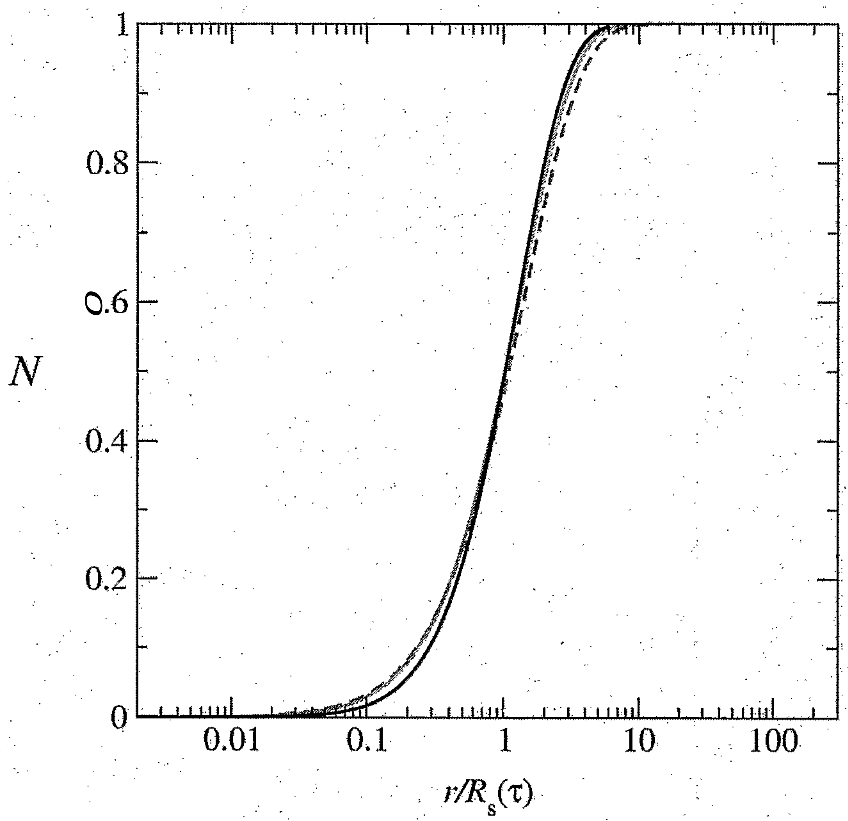


. . . in finite window: lattice artefacts

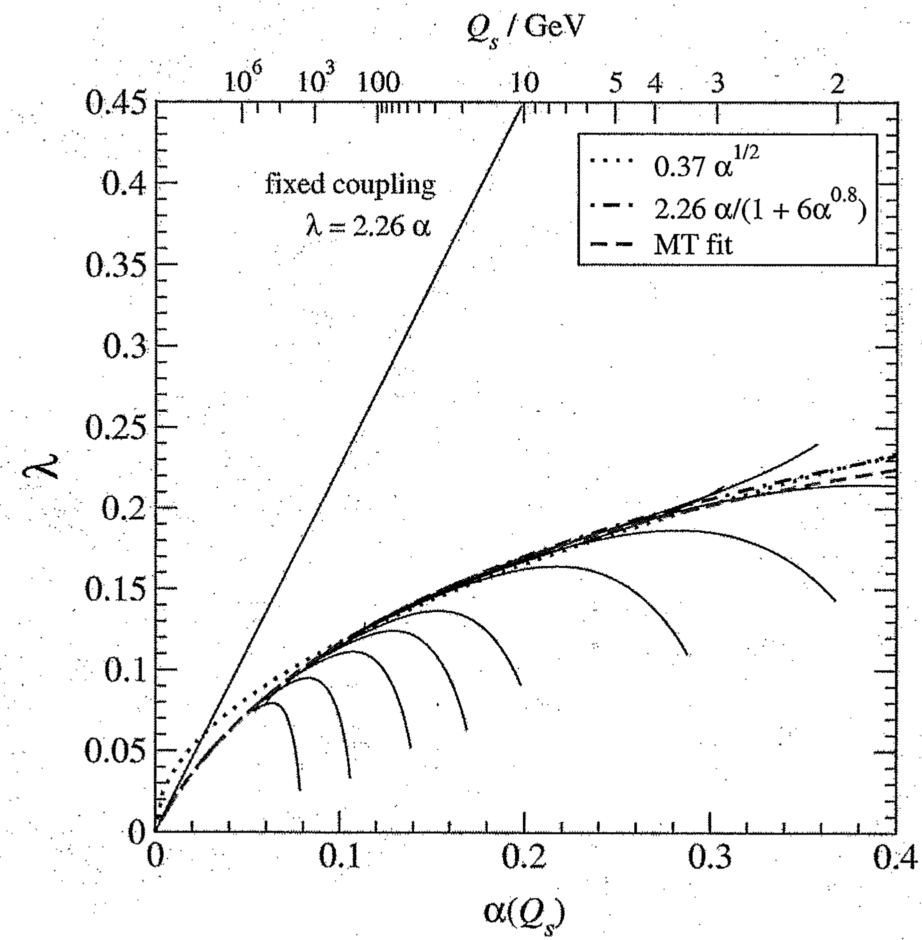
- $Q_s(\eta)$ protects IR ✓
- check UV via $\lambda(\eta) := \partial_\eta \ln Q_s(\eta)$

BK (parent dipole scheme): $\lambda(\eta) = \partial_\eta \ln Q_s(\eta)$

Near scaling despite running coupling



$\lambda(\eta) := \partial_\eta \ln Q_s(\eta)$



▸ Running coupling is essential

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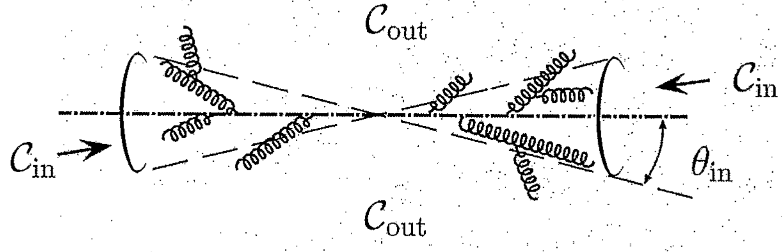
Non-global (exclusive) jet observables

Already seen in $e^+e^- \to$ jets @ total energy E



- fix geometry
- measure soft rad. into $\mathcal{C}_{\text{out}}$ only
- require $\sum E_{\text{soft}} < E_{\text{out}}$
- evolution equation in $\ln(E/E_{\text{out}})$

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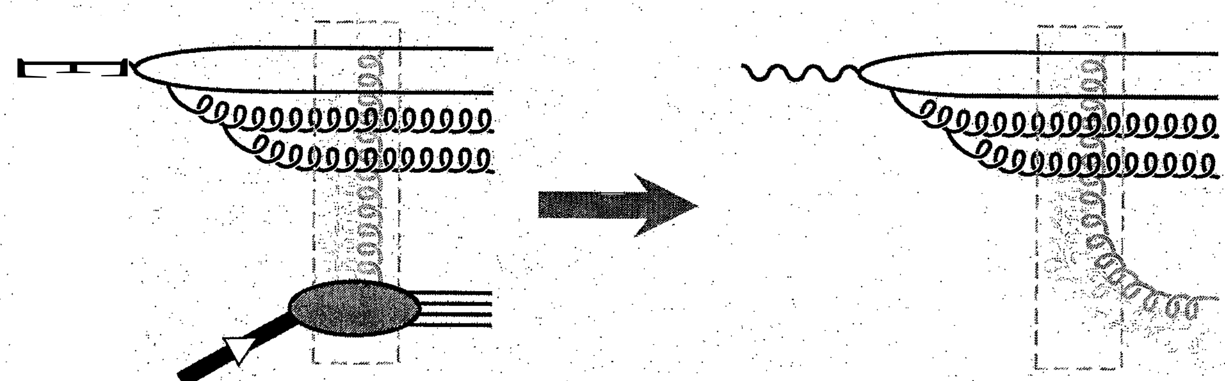
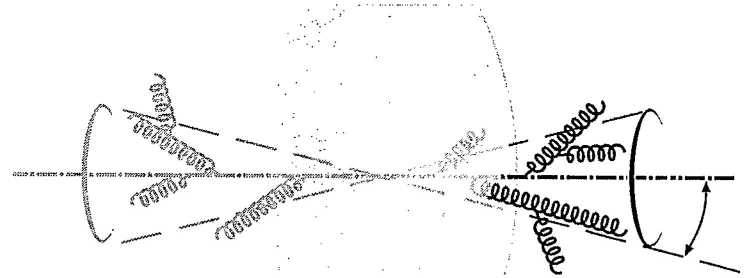
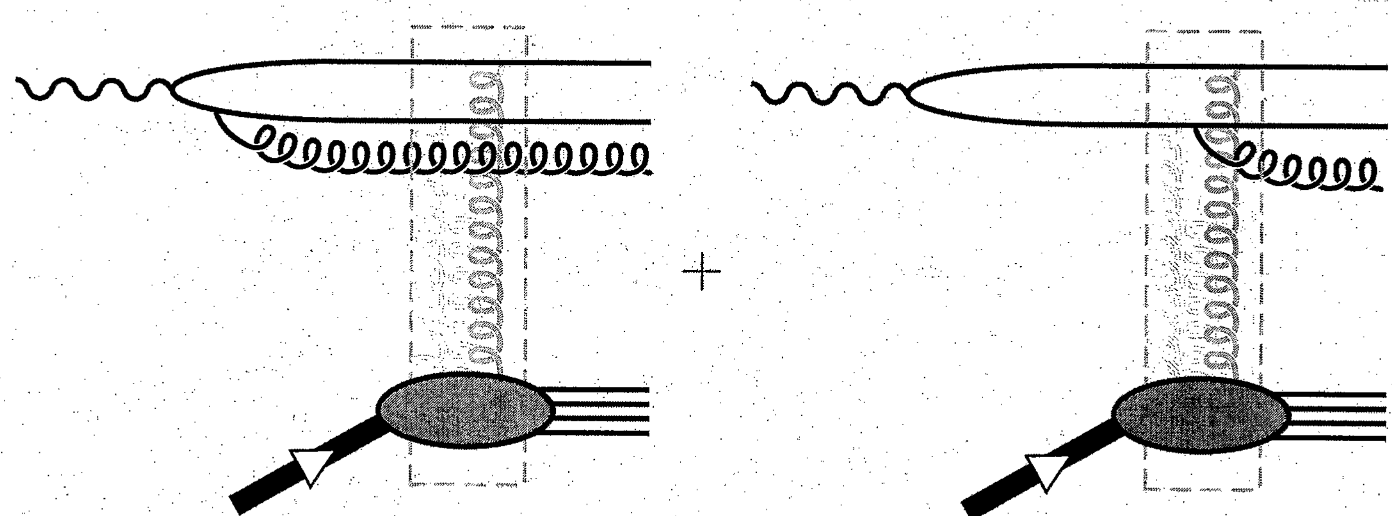
Analogy with CGC amplitudes:



Jets in a medium



Already for global observables: difference to small x



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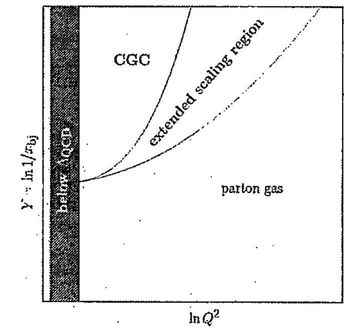
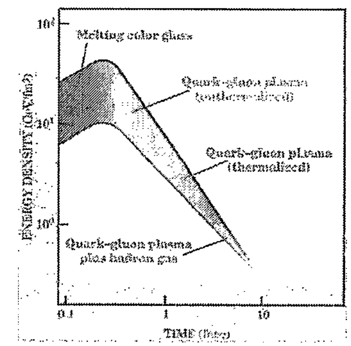
Outlook: tractable problems beyond leading twist

- uses of soft gluon radiation ✓
- CGC in γA ✓
  - saturation scale ✓
  - geometric scaling ✓



- CGC in heavy ion collisions (RHIC & LHC):
  - scales in initial conditions
  - saturation scale & Cronin effect
  - saturation scale & particle multiplicities



- jets in medium @ LHC ✓

in progress

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Perturbative Odderon in the Color Glass Condensate

Yoshitaka Hatta (RBRC)

in collaboration with
E. Iancu, K. Itakura, L. McLerran

We study the perturbative odderon exchange in high energy hadron collisions within the effective theory of Color Glass Condensate. We derive small-x evolution equations for gauge invariant scattering amplitudes describing odderon exchanges between the CGC and two types of projectiles; a color dipole and a system of three quarks. In the linear regime our equations are shown to reduce to the Bartels-Kwiecinski-Praszalowicz equation.

Theoretical status of perturbative odderon

- Bartels-Kwiecinski-Praszalowicz (BKP) equation (1980)

- Mapping onto exactly solvable 1D Heisenberg spin

  Lipatov (1993), Faddeev & Korchemsky (1994)

- Two exact solutions

  Janik & Wosiek (1999) $\alpha_{odd} < 1$

  Bartels, Lipatov & Vacca (2000) $\alpha_{odd} = 1$

- Formulation in Mueller's dipole model

  Kovchegov, Szymanowski & Wallon (2004)

- Formulation in the CGC

Simplifying the JIMWLK equation

Evolution equation for a scattering amplitude

$$\frac{\partial}{\partial\tau}\langle T\rangle_\tau = \int_{xy}\left\langle \frac{\delta}{\delta\alpha^a_\tau(x_\perp)}\eta^{ab}(x_\perp,y_\perp)\frac{\delta}{\delta\alpha^b_\tau(y_\perp)}T\right\rangle_\tau$$

JIMWLK kernel

$$\eta^{ab}(x_\perp,y_\perp)=\frac{1}{\pi}\int\frac{d^2z_\perp}{(2\pi)^2}\frac{(x_\perp-z_\perp)\cdot(y_\perp-z_\perp)}{\underline{(x_\perp-z_\perp)^2(z_\perp-y_\perp)^2}}$$
$$\times\left(1-\tilde{V}^\dagger(x_\perp)\tilde{V}(z_\perp)\right)^{fa}\left(1-\tilde{V}^\dagger(x_\perp)\tilde{V}(y_\perp)\right)^{fb}$$

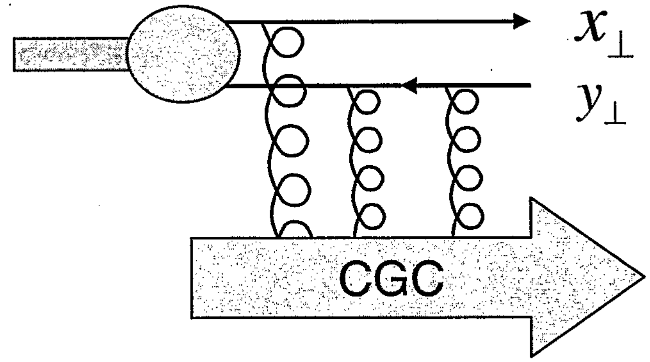
For a gauge invariant amplitude T, the JIMWLK equation can equivalently be written in a manifestly IR finite form.

$$\frac{\partial}{\partial\tau}\langle T\rangle_\tau = -\frac{1}{16\pi^3}\int_{xyz}\frac{(x_\perp-y_\perp)^2}{\underline{(x_\perp-z_\perp)^2(z_\perp-y_\perp)^2}}\left\langle\left(1+\tilde{V}_x^\dagger\tilde{V}_y-\tilde{V}_x^\dagger\tilde{V}_z-\tilde{V}_z^\dagger\tilde{V}_y\right)^{ab}\frac{\delta}{\delta\alpha^a_\tau(x_\perp)}\frac{\delta}{\delta\alpha^b_\tau(y_\perp)}T\right\rangle_\tau$$

"Dipole"- JIMWLK

Construction of the Odderon exchange amplitude in CGC

Dipole-CGC scattering



$$S^{\mathrm{odd}}(x_\perp, y_\perp) = \langle \mathrm{out,\ odd} \,|\, \mathrm{in,\ even} \rangle$$

$$= \frac{1}{2N_c} \left\langle \mathrm{tr}(V_x^\dagger V_y) - \mathrm{tr}(V_y^\dagger V_x) \right\rangle_\tau$$

$$S_{xy}[\alpha] = 1 - N_{xy}[\alpha] + \underbrace{iO_{xy}[\alpha]}_{\text{Odderon amplitude}}$$

Weak field approximation

$$O(x_\perp, y_\perp) \simeq \frac{-g^3}{24N_c} d^{abc} \left\{ 3(\alpha_x^a \alpha_y^b \alpha_y^c - \alpha_x^a \alpha_x^b \alpha_y^c) + (\alpha_x^a \alpha_x^b \alpha_x^c - \alpha_y^a \alpha_y^b \alpha_y^c) \right\}$$

"CGC Green's function"

(dipole) JIMWLK ↓

$$\frac{\partial}{\partial \tau} \langle O(x_\perp, y_\perp) \rangle_\tau = \frac{\bar{\alpha}_s}{2\pi} \int d^2z \frac{(x_\perp - y_\perp)^2}{(x_\perp - z_\perp)^2 (y_\perp - z_\perp)^2} \Big(\langle O(x_\perp, z_\perp) \rangle_\tau + \langle O(y_\perp, z_\perp) \rangle_\tau - \langle O(x_\perp, y_\perp) \rangle_\tau \Big),$$

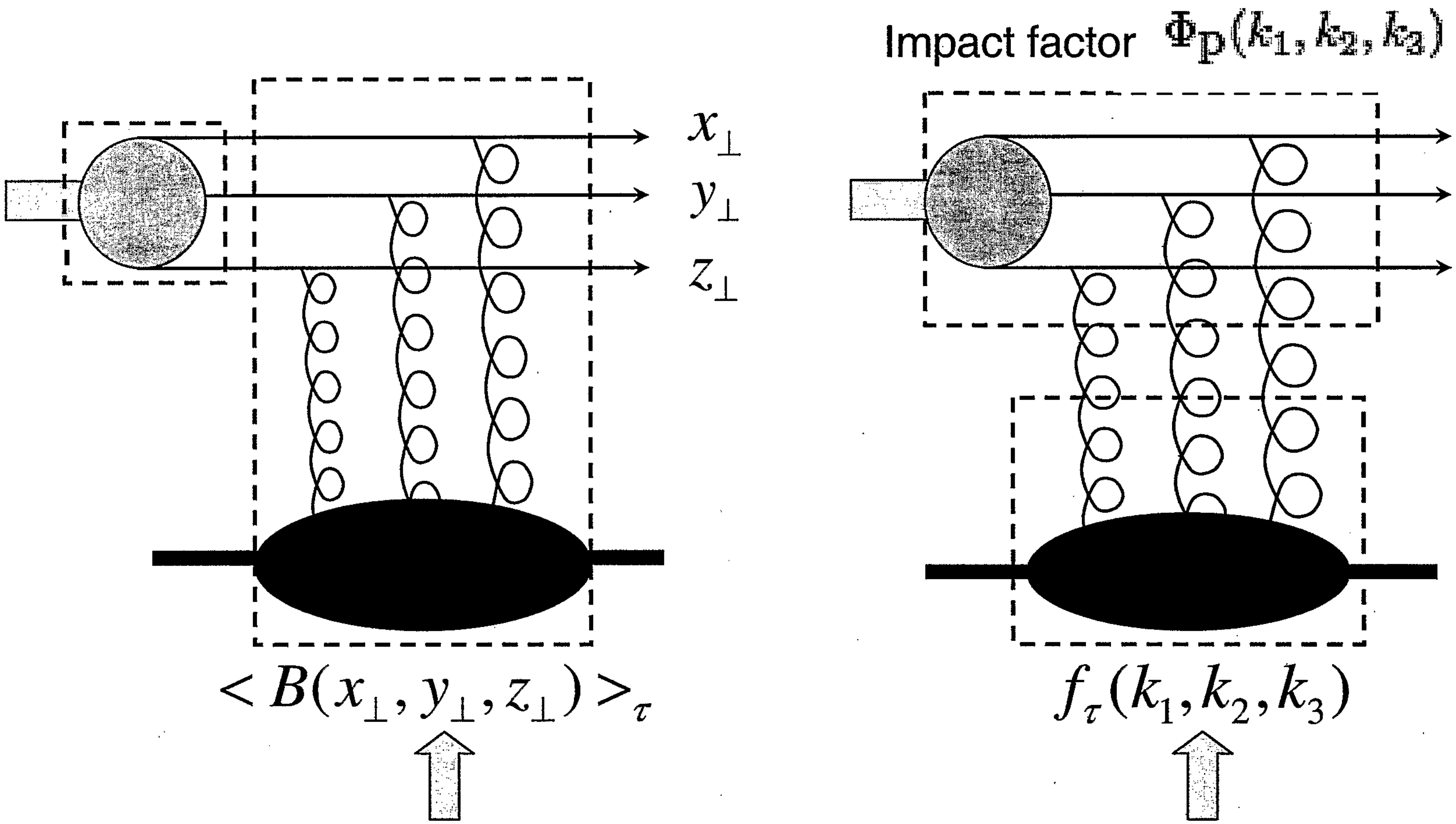
$$\alpha_{odd} = 1$$

Evolution equation for the 3-quark odderon amplitude in the linear regime

$$\frac{\partial}{\partial\tau}\langle B_{xyz}\rangle_\tau = \frac{3\alpha_s}{4\pi^2}\int d^2w_\perp \frac{(x_\perp - y_\perp)^2}{(x_\perp - w_\perp)^2(y_\perp - w_\perp)^2}$$
$$\times\Big(\langle B_{xwz}\rangle_\tau + \langle B_{wyz}\rangle_\tau - \langle B_{xyz}\rangle_\tau$$
$$-\langle B_{wwz}\rangle_\tau - \langle B_{xxw}\rangle_\tau - \langle B_{yyw}\rangle_\tau - \langle B_{xyw}\rangle_\tau\Big)$$
$$+\,(2 \text{ cyclic permutations}).$$

Closed equation for the gauge invariant amplitude $< B(x_\perp, y_\perp, z_\perp) >_\tau$
IR and UV safe
Relation to the BKP equation ?

Equivalence to the BKP equation



JIMWLK equation   BKP equation

Identify $\quad d^{abc} < \alpha_x^a \alpha_y^b \alpha_z^c >_\tau \equiv f_\tau(x_\perp, y_\perp, z_\perp)$

They satisfy the same equation provided one uses the dipole JIMWLK equation for $d^{abc} < \alpha^a \alpha^b \alpha^c >$

The Extended Evolution Equation of the Color Glass Condensate in a Dilute Gluon Medium

S.M.H. Wong

with A.H. Mueller and A. Shoshi

Columbia University, New York

Abstract

The Balitsky-JIMWLK equations are shown to be incomplete as equations that describe the small-x evolution towards the unitarity limit in high energy collisions. Two equivalent approaches to correct this problem are discussed. The JIMWLK equation is extended by a new fourth order functional derivative term and at any given level of the Balitsky hierarchy of equations, each level is now coupled both to the level above and the one below whereas in the original hierarchy, this coupling has always been towards the upward direction only.

Including some parallel work by:

E. Iancu & D. Triantafyllspoulos

[redacted]

- Rare configuration can be more important than mean/typical configuration from BK!
- BK allows unitarity violation in intermediate evolutions $\Longrightarrow$ *something's missing!*
- The JIMWLK equation was thought to describe small-z evolution completely but then numerical simulations showed that there was little difference between BK and JIMWLK!

[Rummukainen & Weigert, NPA **739** p183]

- With gauge invariant operator $\mathcal{O}$

$$\frac{\partial \langle \mathcal{O} \rangle_Y}{\partial Y} = -\frac{1}{16\pi^3} \Big\langle \int_{\boldsymbol{x},\boldsymbol{y}} \mathcal{M}(\boldsymbol{x},\boldsymbol{y},\boldsymbol{z})(1 - \tilde{V}_{\boldsymbol{z}}^\dagger \tilde{V}_{\boldsymbol{x}})^{fa}(1 - \tilde{V}_{\boldsymbol{z}}^\dagger \tilde{V}_{\boldsymbol{y}})^{fb} \times \frac{\delta}{\delta\alpha^a(\boldsymbol{x})}\frac{\delta}{\delta\alpha^b(\boldsymbol{y})}\, \mathcal{O}[\alpha] \Big\rangle_{\boldsymbol{Y}} .$$

$\delta/\delta\alpha$ acts on a Wilson line = a gluon exchange with the Wilson line!

- ◇ JIMWLK and BK have pomeron merging but no splitting! There are no pomeron loops in either! Each $\delta^2/\delta\alpha^a\partial\alpha^a \equiv$ one porneron exchange! Need more $\delta/\delta\alpha$!
- Balitsky hierarchy coupled $\langle T^{(n)} \rangle$'s in only one way — upward!

Finally the extended JIMWLK equation is

$$\frac{\partial}{\partial Y} W_Y[\alpha] = \left(\frac{1}{2} \int_{\boldsymbol{x},\boldsymbol{y}} \frac{\delta}{\delta \alpha_Y^a(\boldsymbol{x})} \eta^{ab}(\boldsymbol{x},\boldsymbol{y}) \frac{\delta}{\delta \alpha_Y^b(\boldsymbol{y})} \right.$$

$$- \frac{g^4}{128\pi^3 N_g} \int_{\substack{\boldsymbol{x},\boldsymbol{y},\boldsymbol{z}\\ \boldsymbol{u_1},\boldsymbol{v_1}\\ \boldsymbol{u_2},\boldsymbol{v_2}}} \mathcal{M}(\boldsymbol{x},\boldsymbol{y},\boldsymbol{z})\, \mathcal{G}(\boldsymbol{u}_1|\boldsymbol{x},\boldsymbol{z}) \mathcal{G}(\boldsymbol{v}_1|\boldsymbol{x},\boldsymbol{z}) \mathcal{G}(\boldsymbol{u}_2|\boldsymbol{z},\boldsymbol{y}) \mathcal{G}(\boldsymbol{v}_2|\boldsymbol{z},\boldsymbol{y})$$

$$\left. \times \frac{6}{\delta\alpha^a(\boldsymbol{u_1})} \frac{\delta}{\delta\alpha^a(\boldsymbol{v_1})} \frac{\delta}{\delta\alpha^b(\boldsymbol{u_2})} \frac{\delta}{\delta\alpha^b(\boldsymbol{v_2})} \nabla_{\boldsymbol{x}}^2 \nabla_{\boldsymbol{y}}^2 \alpha^c(\boldsymbol{x})\, \alpha^c(\boldsymbol{y}) \right) W_Y\, [a].$$

$$= \left(- \frac{\partial}{\partial\alpha^a(\boldsymbol{x})} \sigma^a(\boldsymbol{x}) + \frac{1}{2} \int_{\boldsymbol{x},\boldsymbol{y}} \eta^{ab}(\boldsymbol{x},\boldsymbol{y}) \frac{\delta}{\delta^2 \alpha_Y^a(\boldsymbol{x}) \alpha_Y^b(\boldsymbol{y})} \right.$$

$$- \frac{g^4}{128\pi^3 N_g} \int_{\substack{\boldsymbol{x},\boldsymbol{y},\boldsymbol{z}\\ \boldsymbol{u_1},\boldsymbol{v_1}\\ \boldsymbol{u_2},\boldsymbol{v_2}}} \mathcal{M}(\boldsymbol{x},\boldsymbol{y},\boldsymbol{z})\, \mathcal{G}(\boldsymbol{u}_1|\boldsymbol{x},\boldsymbol{z}) \mathcal{G}(\boldsymbol{v}_1|\boldsymbol{x},\boldsymbol{z}) \mathcal{G}(\boldsymbol{u}_2|\boldsymbol{z},\boldsymbol{y}) \mathcal{G}(\boldsymbol{v}_2|\boldsymbol{z},\boldsymbol{y})$$

$$\left. \times \frac{\delta}{\delta\alpha^a(\boldsymbol{u_1})} \frac{6}{\delta\alpha^a(\boldsymbol{v_1})} \frac{6}{\delta\alpha^b(\boldsymbol{u_2})} \frac{6}{\delta\alpha^b(\boldsymbol{v_2})} \nabla_{\boldsymbol{x}}^2 \nabla_{\boldsymbol{y}}^2 \alpha^c(\boldsymbol{x})\, \alpha^c(\boldsymbol{y}) \right) W_Y[\alpha] \ .$$

where

$$\sigma^a(\boldsymbol{x}) = \frac{1}{2} \int_{\boldsymbol{y}} \frac{\delta}{\delta\alpha^b(\boldsymbol{y})} \eta^{ab}(\boldsymbol{x},\boldsymbol{y}) \ .$$

This extension is necessary:

$$\eta(\boldsymbol{x},\boldsymbol{y}) \sim g^2\alpha^2 + \ldots$$

in the

- Weak field limit (dilute medium): $a \sim$ g

  but not in the

- Strong field limit (dense medium): $a \sim 1/g$

[redacted]

- The JIMWLK equation is a diffusion equation.
  - Brownian motion $\Longrightarrow$ Langevin description.

    [Blaizot, Iancu & Weigert, **NPA 713**, p441]
  - Langevin description $\Longrightarrow$ Computer simulation.
- The extended JIMWLK has fourth order derivative terms and is no longer a diffusion equation.
  - Langevin description ??? Computer simulation ???

Try to recover from these!

The basic form of the extended equation is captured in

$$\frac{\partial}{\partial t}P(x,t) = \Big(-\frac{\partial}{\partial x}a(x) + \frac{\partial^2}{\partial x^2}b(x) - \frac{\partial^4}{\partial x^4}c(x)\Big)P(x,t)\ .$$

$P(x,t)$ is a probability density distribution of some particles whose motions are Markov processes.
Already know that the first two terms can be derived from a drift, $a(x)$, and a Gaussian noise term, $\nu(t)$, in

$$x = a(x) + (2!\,b(x))^{1/2}\nu(t) + (4!\,c(x))^{1/4}\zeta(t)$$

where $\zeta(t)$ is a non-Gaussian noise called a fourth order noise in the literature is required for a stochastic interpretation of the new equation.

$$\frac{\partial}{\partial t} P(x,t) = -\frac{\partial^4}{\partial x^4} c(x)\; P(x,t) \;,$$

$$x = (4!\, c(x))^{1/4}\, \zeta(t) \;.$$

Because of the Markov nature of our model:

- Independent of all past history!
- What is going to happen next depends only on now!
- Bayes' Theorem of conditional probability:

$$P(x,t) = \int \mathrm{d}x' P(x,t|x',t') P(x',t').$$

For a small time increment $At = (t - t') \to 0$,

$$\Delta x(t) = x - x' = \sqrt[4]{4!\; c(x')}\; \zeta(t) \Delta t \;,$$

$$P(x,t|x',t') = \int \mathrm{d}\mathcal{P}_{\mathrm{NG}}(\zeta(t))\, \delta(x - x' - \sqrt[4]{4!\; c(x')}\; \zeta(t)\, \Delta t)$$

Ito's lemma (stochastic calculus) says:

$$P(x,t|x',t') \simeq \left\{ \delta(x - x') - \delta''''(x - x') c(x') \right\} \Delta t + \ldots$$

and

$$P(x',t') = P(x',t) - \frac{\partial}{\partial t} P(x',t) + \ldots .$$

[redacted]

To generate the 4th order diffusion term from a stochastic description, it is necessary to introduce what we called a Non-Diagonal Noise, $\xi(\boldsymbol{x}, \mathrm{y})$

$$\mathrm{d}\mathcal{P}_{\mathrm{ND}}(\xi) = f_{\mathrm{ND}}(\xi)\mathrm{d}\xi$$

$$\langle \mathcal{O}(\xi) \rangle = \int \mathcal{O}(\xi) f_{\mathrm{ND}}(\xi)\mathrm{d}\xi$$

$$\langle \xi(\boldsymbol{x}, \boldsymbol{y}) \rangle_{\mathrm{ND}} = 0$$

$$\boxed{\langle \xi(\boldsymbol{x}_1, \boldsymbol{y}_1)\xi(\boldsymbol{x}_2, \boldsymbol{y}_2) \rangle_{\mathrm{ND}} = \delta^{(2)}(\boldsymbol{x}_1 - \boldsymbol{y}_2)\delta^{(2)}(\boldsymbol{y}_1 - \boldsymbol{x}_2) .}$$

The pairing of the coordinates is "wrong"!
Can't simply take the quartic root of the coefficient function of the $\delta^4/\delta\alpha^4$ terms!

The stochastic equation for α^a is

$$\begin{aligned}\frac{\partial}{\partial Y}\alpha^a(\boldsymbol{u}) &= \sigma^a(\boldsymbol{u}) + \int_{\boldsymbol{z}} \epsilon^{ab,i}(\boldsymbol{u}, \boldsymbol{z})\nu^{b,i}(\boldsymbol{z}) \\ &\quad + \int_{\boldsymbol{x},\boldsymbol{w},\boldsymbol{z}} \rho(\boldsymbol{u}, \boldsymbol{x}, \boldsymbol{w}, \boldsymbol{z})\bar{\nu}^a(\boldsymbol{x}, \boldsymbol{w})\zeta(\boldsymbol{z})\sqrt{\xi(\boldsymbol{x}, \boldsymbol{w})}\end{aligned}$$

$$\rho(\boldsymbol{u}, \boldsymbol{x}, \boldsymbol{w}, \boldsymbol{z}) \propto \frac{1}{\sqrt{(\boldsymbol{x} - \boldsymbol{z})^2}}\{(\boldsymbol{x} - \boldsymbol{w})^2[\nabla^2_{\boldsymbol{x}}\nabla^2_{\boldsymbol{w}}\alpha^c(\boldsymbol{x})\alpha^c(\boldsymbol{w})]\}^{1/4}$$

Thus for two dipoles,

$$\mathcal{O} = T^{(2)}(\boldsymbol{x}_1, \boldsymbol{y}_1; \boldsymbol{x}_2, \boldsymbol{y}_2) = T(\boldsymbol{x}_1, \boldsymbol{y}_1)T(\boldsymbol{x}_2, \boldsymbol{y}_2)$$

$$\begin{aligned}
&\frac{\partial}{\partial Y}\langle T^{(2)}(\boldsymbol{x}_1, \boldsymbol{y}_1; \boldsymbol{x}_2, \boldsymbol{y}_2)\rangle_Y \\
&= \frac{\bar{\alpha}_s}{2\pi}\int_{\boldsymbol{z}}\Bigg(\mathcal{M}(\boldsymbol{x}_1, \boldsymbol{y}_1, \boldsymbol{z}) \otimes \langle T^{(2)}(\boldsymbol{x}_1, \boldsymbol{y}_1; \boldsymbol{x}_2, \boldsymbol{y}_2)\rangle_Y \\
&\qquad -\mathcal{M}(\boldsymbol{x}_1, \boldsymbol{y}_1, \boldsymbol{z})\langle T^{(3)}(\boldsymbol{x}_1, \boldsymbol{z}; \boldsymbol{z}, \boldsymbol{y}_1; \boldsymbol{x}_2, \boldsymbol{y}_2)\rangle_Y\Bigg) \\
&\qquad +(1 \longleftrightarrow 2) \\
&\quad +\mathbf{T^{(1)}}!! \quad \Longleftarrow \quad \textbf{THERE IS A NEW TERM}
\end{aligned}$$

In general for n dipoles,

$$\mathcal{O} = T^{(n)}(\boldsymbol{x}_1, \boldsymbol{y}_1; \ldots; \boldsymbol{x}_n, \boldsymbol{y}_n) = T(\boldsymbol{x}_1, \boldsymbol{y}_1)\ldots T(\boldsymbol{x}_n, \boldsymbol{y}_n)$$

The MODIFIED Balitsky hierarchy

$$\begin{aligned}
&\frac{\partial}{\partial Y}\langle T^{(n)}(\boldsymbol{x}_1, \boldsymbol{y}_1; \ldots; \boldsymbol{x}_n, \boldsymbol{y}_n)\rangle_Y \\
&= \frac{\bar{\alpha}_s}{2\pi}\sum_{i=1}^{n}\int_{\boldsymbol{z}}\Bigg(\mathcal{M}(\boldsymbol{x}_i, \boldsymbol{y}_i, \boldsymbol{z}) \otimes \langle T^{(n)}(\boldsymbol{x}_1, \boldsymbol{y}_1; \ldots; \boldsymbol{x}_i, \boldsymbol{y}_i; \ldots; \boldsymbol{x}_n, \boldsymbol{y}_n)\rangle_Y \\
&\qquad -\mathcal{M}(\boldsymbol{x}_i, \boldsymbol{y}_i, \boldsymbol{z})\langle T^{(n+1)}(\boldsymbol{x}_1, \boldsymbol{y}_1; \ldots; \boldsymbol{x}_i, \boldsymbol{z}; \boldsymbol{z}, \boldsymbol{y}_i; \ldots; \boldsymbol{x}_n, \boldsymbol{y}_n)\rangle_Y\Bigg) \\
&\qquad +\mathbf{T^{(n-1)}}!! \quad \Longleftarrow \quad \textbf{THERE IS A NEW TERM}
\end{aligned}$$

[redacted]

- In a dilute gluon medium, there are $\delta^4/\delta\alpha^4$ terms that are important for small-x evolution of $W_Y[a]$!
- Same term is responsible for the modified Balitsky hierarchy of equations. $\langle T^{(n)}\rangle$ is now linked to both $\langle T^{(n+1)}\rangle$ and $\langle T^{(n-1)}\rangle$ in one equation.
- New terms reproduce the correct $(\text{pomeron})^3$-vertex!
- Problems:
  - ◇ Yet to show that $S(\boldsymbol{x},\boldsymbol{y})$ approaches unitarity at the correct rate! x
  - ◇ Yet to show unitarity can be restored in the intermediate evolutions! x
  - ◇ With $\delta^4/\delta\alpha^4$, a pomeron from the CGC can now split into two! Pomeron loops by iteration! $\surd$
- New problems:
  - ◇ 4th order diffusion spoiled the Langevin description! x
  - ◇ But able to preserve the stochastic description, introduced 4th order noise! $\surd$
  - ◇ Non-diagonal noise, computer simulation ??
- Extended equation is a sum of very different parts: CGC +color dipole $\Longleftarrow$ Can we do better?

J.Bartels, Hamburg University:

Reggeon Field Theory in QCD

Abstract:

In the first part I review the concept of reggeon field theory in QCD. Reggeon field theory provides a solution,to t-channel reggeon unitarity equations. In QCD, the reggeized gluon plays the role of the fundamental field, and there exist momentum dependent vertex functions which describe the interactions between reggeized gluons. Known examples include the BFKL kernel, the $2 \to 4$ and the $2 \to 6$ gluon transition vertex functions. At present the BFKL kernel is known in NLO accuracy, the other vertex functions have been computed in leading order only. With these vertex functions it is possible to compute composite states of two or three reggeized gluons, Pomerons or Odderons, resp. Another example of interest is the Pomeron loop correction to the Pomeron self energy. Contact to the Color Glass Condensate is made by transforming reggeon field theory from momentum space to configuration space: as a first example, it can be shown that the Fourier transform of the $2 \to 4$ gluon transition vertex - with certain approximations - coincides with the kernel of the Balitsky-Kovchegov equation. Reggeon field theory also allows to compute higher order corrections, both in $1/N_c$ and in α_s.

In the second part of the talk I address a few special topics which can be derived from QCD reggeon field theory: this includes the description of diffraction in deep inelastic scattering, and the validity of the AGK cutting rules in pQCD.

Reggeon Field Theory in QCD

J. Bartels
II.Inst.f.Theor.Physik, Univ.Hamburg

CGC Workshop, BNL, March 2005

Content:

- Introduction: motivation
- Foundation:
  reggeon unitarity equations in momentum space
- Results: reggeon vertices
- Reggeization, bootstrap
- Diffractive vertex
- AGK cutting rules
- Conclusions

Introduction: motivation

Aim: how to formulate and derive a field theory of interacting Pomerons in QCD?

Now very polular: JIMWLK

- in configuration space
- s-channel picture (color glass condensate)
- 'semiclassical framework' for Pomeron interactions (BK equation); recently: Pomeron loops
- (large N_c limit)

Complementary: reggeon field theory in QCD

- derived in momentum space
- t-channel approach
- reggeon unitarity equations, reggeon field theory as solution to these unitarity equations
- allows to compute corrections in α_s (and in $1/N_c^2$).

This talk: second approach, con plementary to most of the talks at this meeting.
Connection with cross sections, co linear hard scattering.

Current status.

Outlook

What has been accomplished:

- formulation of reggeon field theory in QCD
- some NLO elements have been computed, both in momentum space and in configuration space
- framework for NLO caculations
- understand AGK rules in pQCD

What needs to be done:

- How to find solutions?
- Before that: lots to compute in momentum space, e.g. inclusive cross section formulae in heavy ion collisions
- understand connection between reggeon field theory and conformal field theory
  (Virasoro algebra, halomorphic separability, conformal bootstrap, AdS/CFT correspondence,...)

B-JIMWLK^2: beyond & beside.

A. Kovner (UConn)

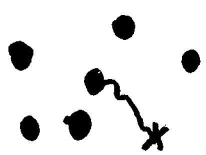
I discuss the derivation of the evolution equation for a dilute projectile. I show how path ordered exponentials of the derivative with respect to colour charge density of the target arises. The meaning of this exponential as the eikonal scattering amplitude of an arbitrary projectile on a single gluon of the target is explained. I also derive the property of self duality of the high energy evolution kernel valid in eikonal approximation.

A.K. & M. Lublinsky, hep-ph/0501198
hep-ph/0502071
hep-ph/0502119

What is missing?

- **Bleaching of color**

Gluons emitted on top of other gluons:



- emitted into a spot that is taken already.

If ρ large, it does not grow linearly but random walks

Maybe expect $y \to \sqrt{y}$?

- **Emission is collective.**

$$b_i \text{ "}\sim\text{" } \frac{D_i}{D^2}\rho \ll \frac{\partial_i}{\partial^2}\rho$$

Are these Pomeron loops?

Some comments.

χ^{JIMWLK}:

$$2\frac{\delta}{\delta\alpha(x)}\frac{\delta}{\delta\alpha(y)} - 2\frac{\delta}{\delta\alpha(x)}\underbrace{Pe^{i\int dx^- T^a\alpha^a(z)}}\frac{\delta}{\delta\alpha(y)}$$

"gluon of the projectile

χ^{KLWMIJ}

$$2\rho(x)\rho(y) - 2\rho(x)\underbrace{Pe^{\int dx^- T^a\frac{\delta}{\delta\rho^a(z)}}}\rho(y)$$

"gluon of the target.

eikonal amplitude of the scattering of the whole projectile on one gluon of the target.

<u>One</u>-<u>gluon</u> <u>projectile</u>:

$$S(x) \simeq 1 + i\, T^a \alpha^a(x)$$

$$\int d\rho\, \underbrace{S(x)}_{P\exp\left[\int dx^- T\cdot\frac{\delta}{\delta\rho}\right]} R(y) W[\rho] = 1 + ig\, \underbrace{\frac{1}{\partial^2}(x-y)}_{\text{one gluon exchange}}$$

<u>Two</u>-<u>gluon</u> <u>projectile</u>:

$$\int d\rho\, S(x) S(y)\, R(z)\, W[\rho] =$$

$$= 1 + ig\left[\frac{1}{\partial^2}(x-z) + \frac{1}{\partial^2}(y-z)\right] - \frac{g^2}{2}\, \underbrace{\frac{1}{\partial^2}(x-z)\, \frac{1}{\partial^2}(y-z)}_{\text{both incoming gluons scatter}}$$

ETC

Many derivatives in R are important if there are many impinging partons - Large projectile.

χ^{KLWMIJ} resumms unitarization corrections due to scattering of many projectile partons on a single parton of the target.

Ω's fluctuations - almost nothing in the target most of the time, even an elephant will not do anything.

Some techical advances:

Can get rid explicitly of x^-.

Can derive dipole limit :

$\mathrm{tr}\, R^\dagger(x) R(y)$ dipole of the target.

Expand to $O\left(\frac{\delta^4}{\delta\rho^4}\right)$ - result does not coincide with MSW -

- can see that 2 dipoles scatter on 1 - gives contribution.

<u>Next ?</u>

Can we figure out Lorentz transformation exactly?

d

Quark production in pA collisions: rescatterings and kt-factorization breaking

François Gelis

We model proton-nucleus collisions in the Color Glass Condensate framework by assuming that the proton can be described by a weak color source that we treat at leading order, while the nucleus is described by a strong color source that needs to treated to all orders. At this level of approximation, the classical Yang-Mills equations can be solved in closed form.

One can then calculate the propagator of a quark in this background field in order to obtain the amplitude for quark-antiquark pair production in proton-nucleus collisions. By squaring this amplitude, one finds that the quark production cross-section involves correlators of 2 and **3** Wilson lines. The fact that the cross-section contains this 3-point correlators implies that one cannot write it in $k_\perp$-factorized form: one recovers $k_\perp$-factorization only in certain limits for which the final state contains some scale which is much larger than the saturation momentum in the nucleus. Numerically, we find the following patterns for the breaking of kl-factorization:

- The magnitude of the breaking of $k_\perp$-factorization decreases as the quark mass increases. Indeed, since the terms that break $k_\perp$-factorization correspond to extra rescatterings, it is natural that massive quarks are less sensitive to these effects than light quarks.

- The magnitude of the breaking of $k_\perp$-factorization is maximal for a transverse momentum $q_\perp \approx Q_s$ of the quark, where Q_s is the saturation scale in the nucleus. One recovers $k_\perp$-factorization when the quark transverse momentum becomes much larger than all the other scales.

- .If Q_s remains smaller or comparable to the quark mass and transverse momentum, the corrections due to the breaking of $k_\perp$-factorization enhance the cross-section. This is interpreted as a threshold effect: having more rescatterings tend to push a few more $Q\overline{Q}$ pairs just above the kinematical production threshold.

- If Q_s is large compared to the mass of the quark, then the corrections due to the breaking of $k_\perp$-factorization tend to reduce the cross-section at small transverse momentum. Since the typical momentum transfer in a scattering is of the order of Q_s, it is indeed more difficult to produce light quarks with a small momentum if they scatter more.

Quark production in pA collisions: rescatterings and kt-factorization breaking

François Gelis

CEA / DSM / SPhT

Pair production amplitude

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- Total amplitude:

$$\mathcal{M}_F = g^2 \int\limits_{\vec{\boldsymbol{k}}_{1\perp},\vec{\boldsymbol{k}}_\perp} \frac{\rho_{p,a}(\vec{\boldsymbol{k}}_{1\perp})}{k_{1\perp}^2} \int\limits_{\vec{\boldsymbol{x}}_\perp,\vec{\boldsymbol{y}}_\perp} e^{i\vec{\boldsymbol{k}}_\perp\cdot\vec{\boldsymbol{x}}_\perp} e^{i(\vec{\boldsymbol{p}}_\perp+\vec{\boldsymbol{q}}_\perp-\vec{\boldsymbol{k}}_\perp-\vec{\boldsymbol{k}}_{1\perp})\cdot\vec{\boldsymbol{y}}_\perp}$$

$$\times\overline{u}(\vec{\boldsymbol{q}})\Big\{[\widetilde{U}(\vec{\boldsymbol{x}}_\perp)t^a\widetilde{U}^\dagger(\vec{\boldsymbol{y}}_\perp)]T_{q\bar{q}}(\vec{\boldsymbol{k}}_\perp) + [t^b U_{ba}(\vec{\boldsymbol{x}}_\perp)]\not{L}\Big\}v(\vec{\boldsymbol{p}})$$

with

$$T_{q\bar{q}}(\vec{\boldsymbol{k}}_\perp) \equiv \frac{\gamma^+(\not{q}-\not{k}+m)\gamma^-(\not{q}-\not{k}-\not{k}_1+m)\gamma^+}{2p^+[(\vec{\boldsymbol{q}}_\perp-\vec{\boldsymbol{k}}_\perp)^2+m^2]+2q^+[(\vec{\boldsymbol{q}}_\perp-\vec{\boldsymbol{k}}_\perp-\vec{\boldsymbol{k}}_{1\perp})^2+m^2]}$$

- **Notes:**
  - the V's cancel between regular and singular contributions
  - $\overline{u}(\vec{\boldsymbol{q}})\left[\not{C}_U(p+q,\vec{\boldsymbol{k}}_{1\perp}) - \gamma^+\frac{(p+q)^2}{p^++q^+}\right]v(\vec{\boldsymbol{p}}) = \overline{u}(\vec{\boldsymbol{q}})\not{L}v(\vec{\boldsymbol{p}})$

Single quark cross-section

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- Single quark production cross-section:

$$\frac{d\sigma_q}{d^2\vec{\boldsymbol{q}}_\perp dy_q} = \frac{\alpha_s^2 N}{8\pi^4 d_A}\int\frac{dp^+}{p^+}\int_{\vec{\boldsymbol{k}}_{1\perp},\vec{\boldsymbol{k}}_{2\perp}} \frac{1}{\boldsymbol{k}_{1\perp}^2\boldsymbol{k}_{2\perp}^2}$$

$$\times\Big\{\mathrm{tr}\Big[(\not{q}+m)T_{q\bar{q}}(\vec{\boldsymbol{k}}_{2\perp})(\not{p}-m)T^*_{q\bar{q}}(\vec{\boldsymbol{k}}_{2\perp})\Big]\frac{C_F}{N}\phi_A^{q,q}(\vec{\boldsymbol{k}}_{2\perp})$$

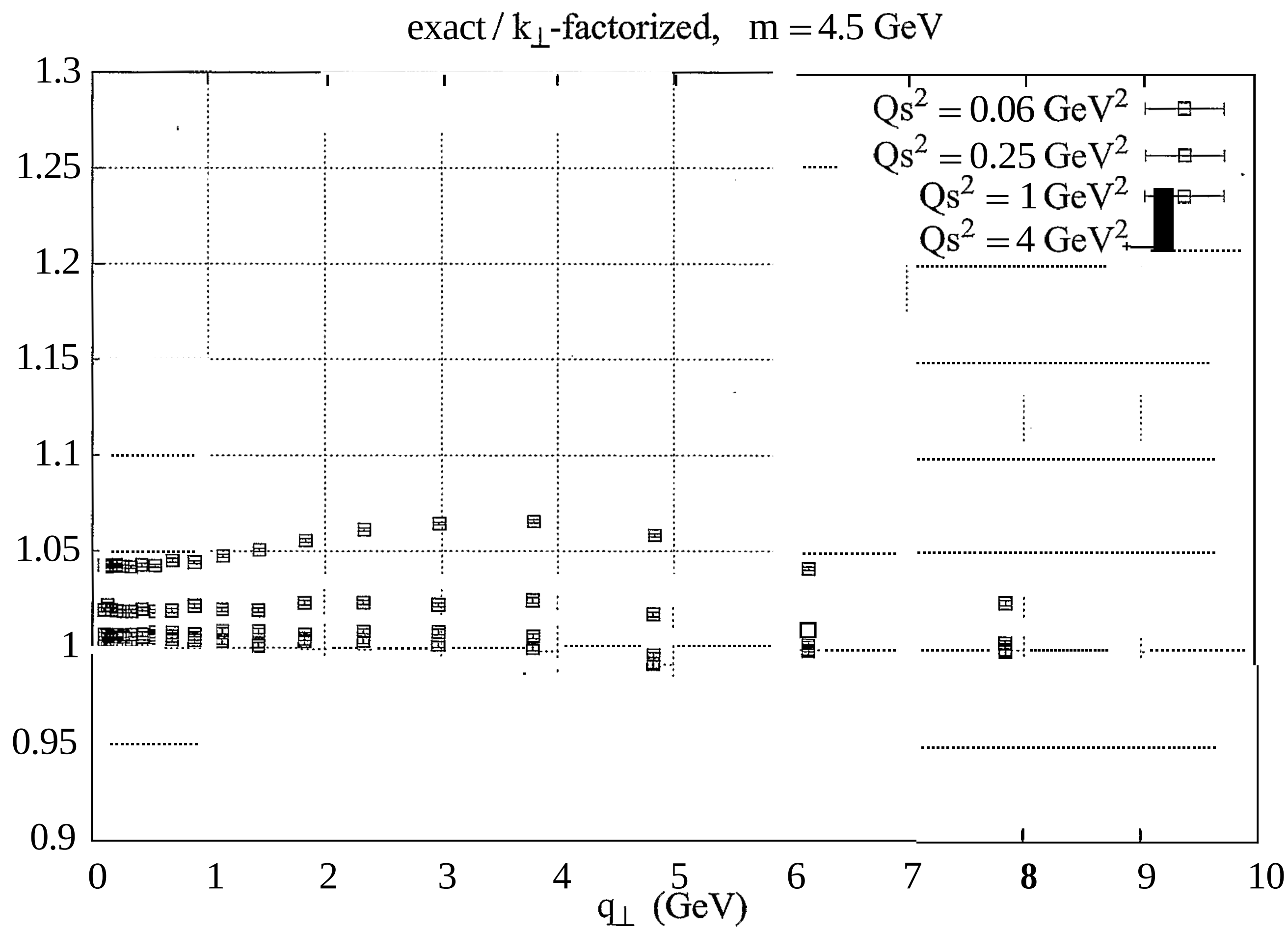
$$+\int_{\vec{\boldsymbol{k}}_\perp}\mathrm{tr}\Big[(\not{q}+m)T_{q\bar{q}}(\vec{\boldsymbol{k}}_\perp)(\not{p}-m)\not{L}^* + \mathrm{h.c.}\Big]\phi_A^{q\bar{q},g}(\vec{\boldsymbol{k}}_{2\perp}|\vec{\boldsymbol{k}}_\perp)$$

$$+\mathrm{tr}\Big[(\not{q}+m)\not{L}(\not{p}-m)\not{L}^*\Big]\phi_A^{g,g}(\vec{\boldsymbol{k}}_{2\perp})\Big\}\varphi_p(\vec{\boldsymbol{k}}_{1\perp})$$

  - $\phi_{\boldsymbol{A}}^{q,q}$ is the analogue of $\phi_A^{g,g}$ for the fundamental representation
  - $k_\perp$-factorization still broken for the nucleus
  - contains only 2-point and 3-point correlators

Breaking of Kt factorization

- Single b-quark cross-section (large N):



- Pair production amplitude
- Single quark cross-section
- Breaking of Kt factorization
- Breaking of Kt factorization
- Breaking of Kt factorization

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Breaking of Kt factorization

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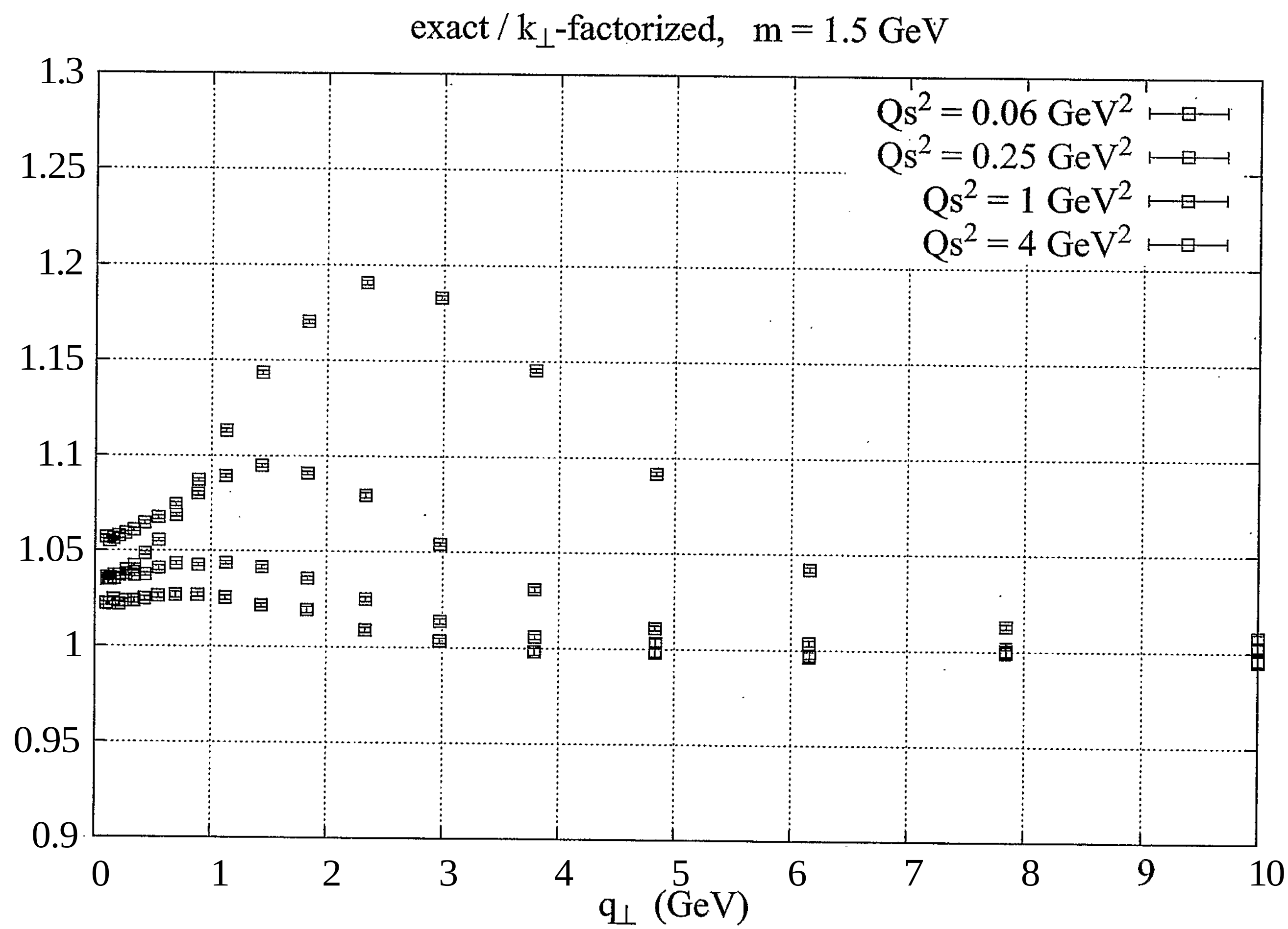
- Single c-quark cross-section (large N) :



Breaking of Kt factorization

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- General trends for the breaking of $k_\perp$-factorization :
  - The magnitude of the breaking increases as m decreases
  - The magnitude of the breaking increases with Q_s
  - The effect is maximum for $q_\perp \sim Q_s$
  - $k_\perp$-factorized is recovered at large $q_\perp$
  - If $Q_s \lesssim m, q_\perp$, the $k_\perp$-factorization breaking terms enhance the cross-section: having more scatterings pushes a few more pairs above the kinematical threshold
  - If $Q_s \gg m, q_\perp$, the effect is a reduction of the cross-section: with a large Q_s it becomes less likely to produce a quark with a small transverse mass

  ⇒ These corrections tend to enhance the Cronin peak that one would obtain by using the $k_\perp$-factorized formula for quark production

Geometric scaling, behavior of the saturation scale and experimental data

Carlos A. Salgado

CERN Physics Department, Theory Division, CH-1211 Geneva

We solve numerically the Balitsky–Kovchegov equations for the cases of running and fixed strong coupling constant. In agreement with previous numerical and analytical results, we find that an asymptotic scaling solution appears for both cases. The small-r behavior of these solutions is, however, different: an anomalous dimension of $\gamma \sim 0.65$, in agreement with analytical estimations is found for fixed coupling; in contrast, $\gamma \simeq 0.85$ for the running coupling case. The rapidity and nuclear size dependence of the saturation scale are also computed and found to be in good agreement with analytical estimations.

Lepton-proton experimental data are known to present *geometric scaling.* We generalize this scaling to the nuclear case and found, in this way, that the A-dependence of the saturation scale is faster than $A^{1/3}$. This geometric scaling is then used to describe the multiplicities measured in nucleus-nucleus and proton-proton collisions at central rapidities for which a very simple formula is derived. The suppression of particles with high-p_t at the forward rapidity region of RHIC are also described by the same geometric scaling. The possible relations of these findings with the numerical solutions of the BK equations are commented.

Based on:

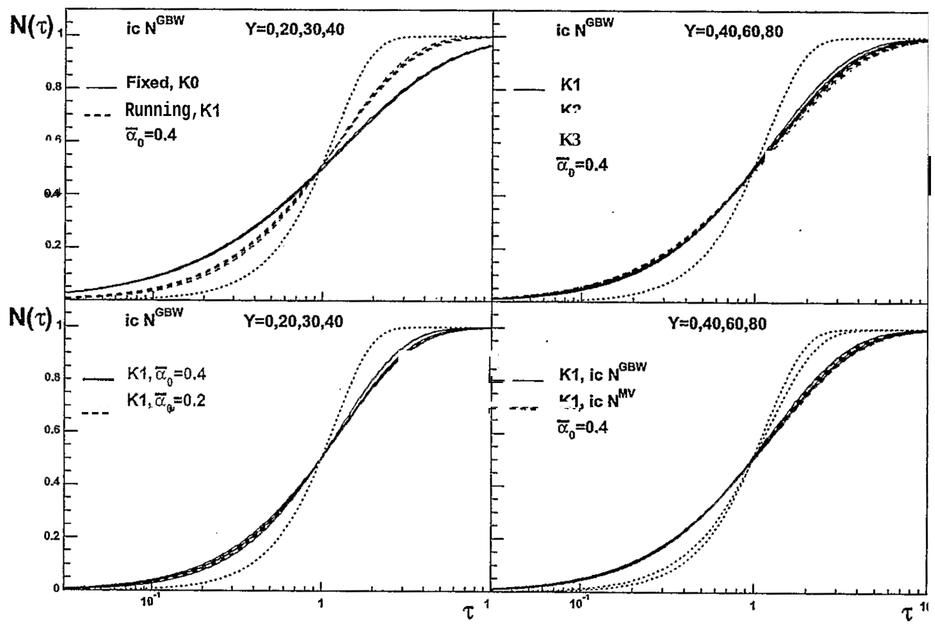
Albacete, Armesto, Kovner, Salgado, Wiedemann PRL92 (2004) 082001

Albacete, Armesto, Milhano, Salgado, Wiedemann PRD71 (2005) 014003

Armesto, Salgado, Wiedemann PRL94 (2005) 022002

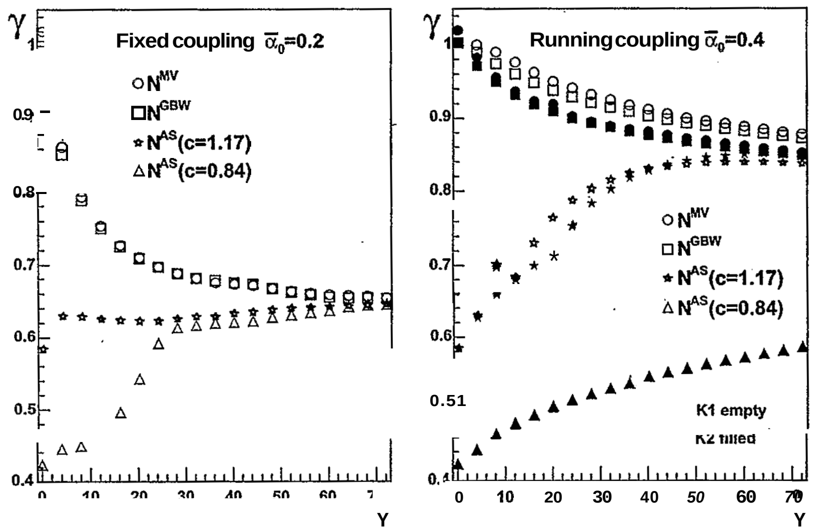
Scaling

[Albacete, Armesto, Milhano, Salgado, Wiedemann 2004]



- $\tau \equiv rQ_{\rm sat}(Y) \Longrightarrow N(\tau)$
- We define $Q_{\rm sat}(Y)$ so that $N(r=1/Q_{\rm sat}(Y),Y) = \kappa$; $\kappa = 1/2$
- Scaling also for the running coupling case

Scaling function



- Fits to $N = a\tau^{2\gamma}(\log\tau + \delta)$ in the region $10^{-5} < \tau < 10^{-1}$
  - Fixed coupling value $\gamma \sim 0.65$ agrees with analytical $\gamma \sim 0.63$
  - Running coupling value $\gamma \sim 0.85$ <u>different</u> from fixed

Geometric scaling in lepton-nucleus data

- Scaling when

$$\frac{\sigma^{\gamma^* A}(\tau)}{\pi R_A^2} = \frac{\sigma^{\gamma^* p}(\tau)}{\pi R_p^2} .$$

- We define

$$Q^2_{\rm sat,A} = Q^2_{\rm sat,p} \left(\frac{A R_p^2}{R_A^2} \right)^{1/\delta}$$

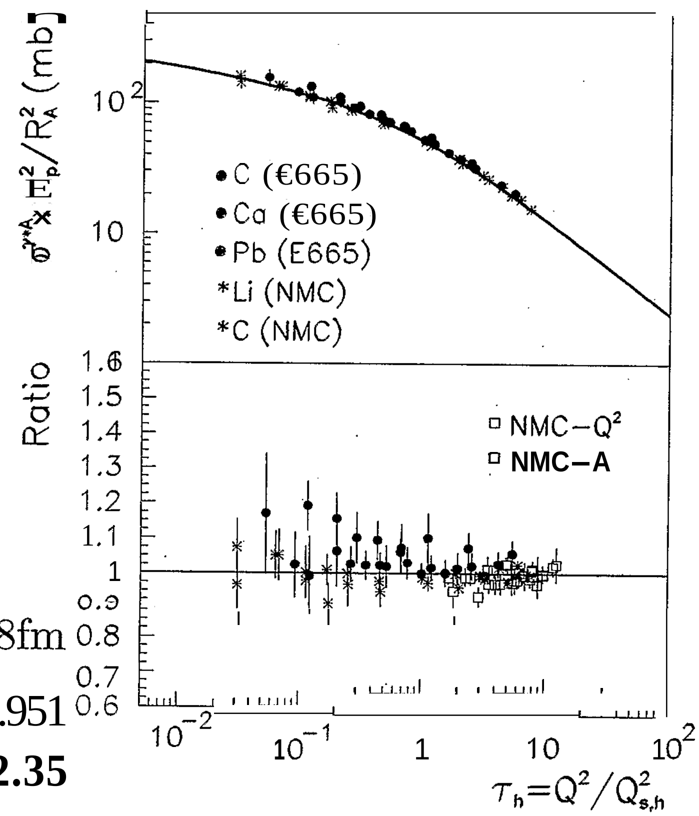
$$R_A = 1.12\, A^{1/3} - 0.86\, A^{-1/3}$$

- R_p, δ free parameters

$$\boxed{\delta = 0.79 \pm 0.02}\; R_p = 0.70 \pm 0.08 \mathrm{fm}$$

$$\Rightarrow Q^2_{\rm sat} \sim A^{4/9} \quad [\chi^2/\mathrm{dof} = 0.951$$

$$\delta = 1 [Q^2_{\rm sat} \sim A^{1/3}] \Rightarrow \chi^2/\mathrm{dof} = \mathbf{2.35}$$



[Armesto, Salgado; Wiedemann (2005)]

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Multiplicities and geometric scaling

⇛ The multiplicit es

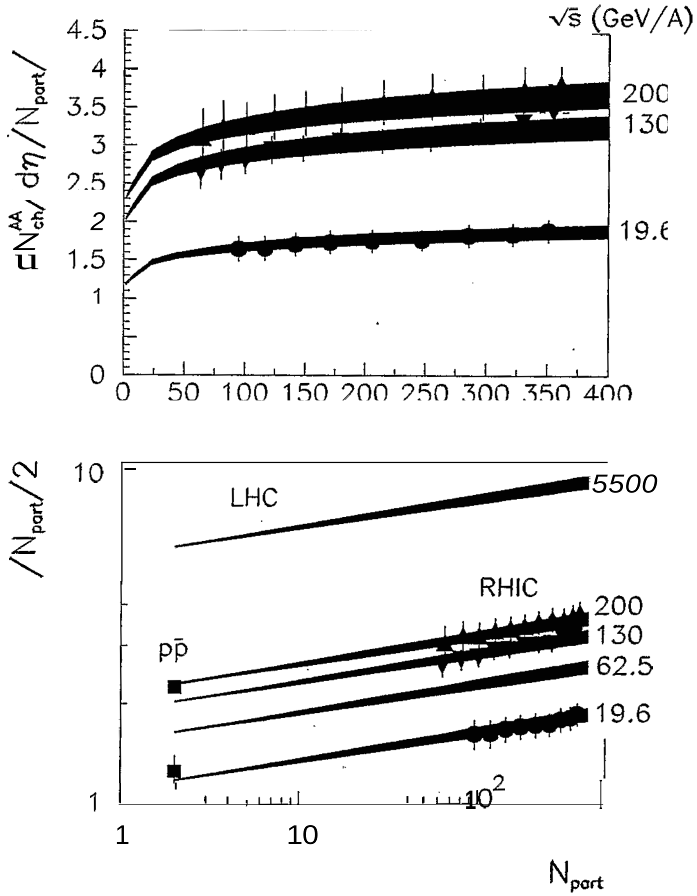
$$\frac{1}{N_{\rm part}}\frac{dN^{AA}}{d\eta}\bigg| = N_0\sqrt{s}^{\chi} N_{\rm part}^{\frac{1-\delta}{3\delta}} .$$

(lepton-proton data)

⇛

⇛

Data: PHOBOS PRC65, 061901 (2002); nucl-ex/0405027



[Armesto, Salgado, Wiedemann (2005)]

Conclusions

- Solution of BK found numerically for running and fixed coupling
  - Asymptotic scaling curve
  - Small-r behavior agrees with analytical results – however, different γ for running and fixed coupling
  - DLL found but no clear scaling window could be identified
- $Q^2_{\rm sat}(y;A)$ agrees with analytical estimates for fixed and running α_S
- The geometric scaling found in lp data has been extended to the nuclear case
  - $Q^2_{\rm sat,A}$ grows faster than $A^{1/3}$.
- Nice description of multiplicities in AA at y $=0$ and suppression of particle production at forward rapidities.
- Numerical coincidence, saturation?? $\longrightarrow$ use evolution eqs.
- DGLAP nuclear gluons are constrained for $x \gtrsim 0.02$ by DIS data
  - Check universality of nPDF with RHIC data

Ian Balitsky

Scattering of shock waves **in QCD.**

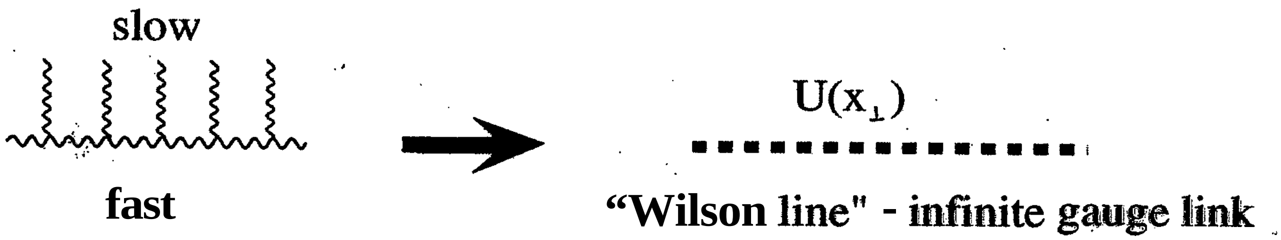
Viewed from the center of mass frame, a typical high-energy hadron-hadron scattering looks like a collision of two shock waves. Indeed, due to the Lorentz contraction the two hadrons shrink into thin "pancakes" which collide producing the final state particles. The mail question is the field/particles produced by the collision of these two waves. On the theoretical side, this is related to the problem of high-energy effective action and to the the ultimate question of the small-x physics - unitarization of the BFKL pomeron and the Froissart bound in QCD[?]. On more practical terms, the immediate result of the scattering of the two shock waves gives the initial conditions for the formation of a quark-gluon plasma observed in the heavy-ion collisions at RHIC.

Due to parton saturation at high energies, the collision of QCD shock waves can be studied using semiclassical methods.

At present, the corresponding Yang-Mills equations has not been solved analytically. In my talk I formulate the problem of scattering of shock waves, find the boundary conditions for the double functional integral for the cross section, develop the expansion in the commutators of two shock waves equivalent to the series in strength of one of the waves, and calculate the second-order term of this expansion.

Key observation:

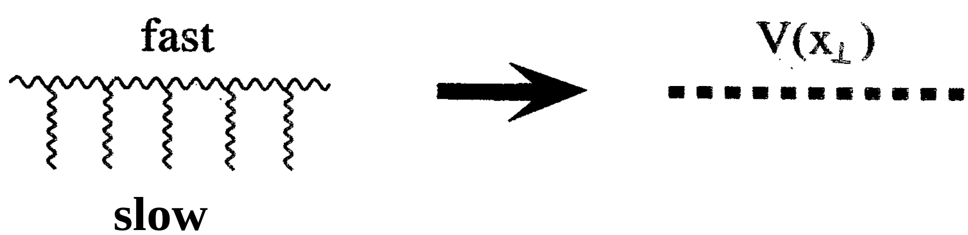
In the frame of the spectator



$$U(x_\perp) = [-\infty n_1 + x_\perp, \infty n_1 + x_\perp]$$

$$[x, y] \equiv P e^{ig\int_0^1 dt (x-y)^\mu A_\mu(tx+(1-t)y)}$$

Similarly, in the lab frame



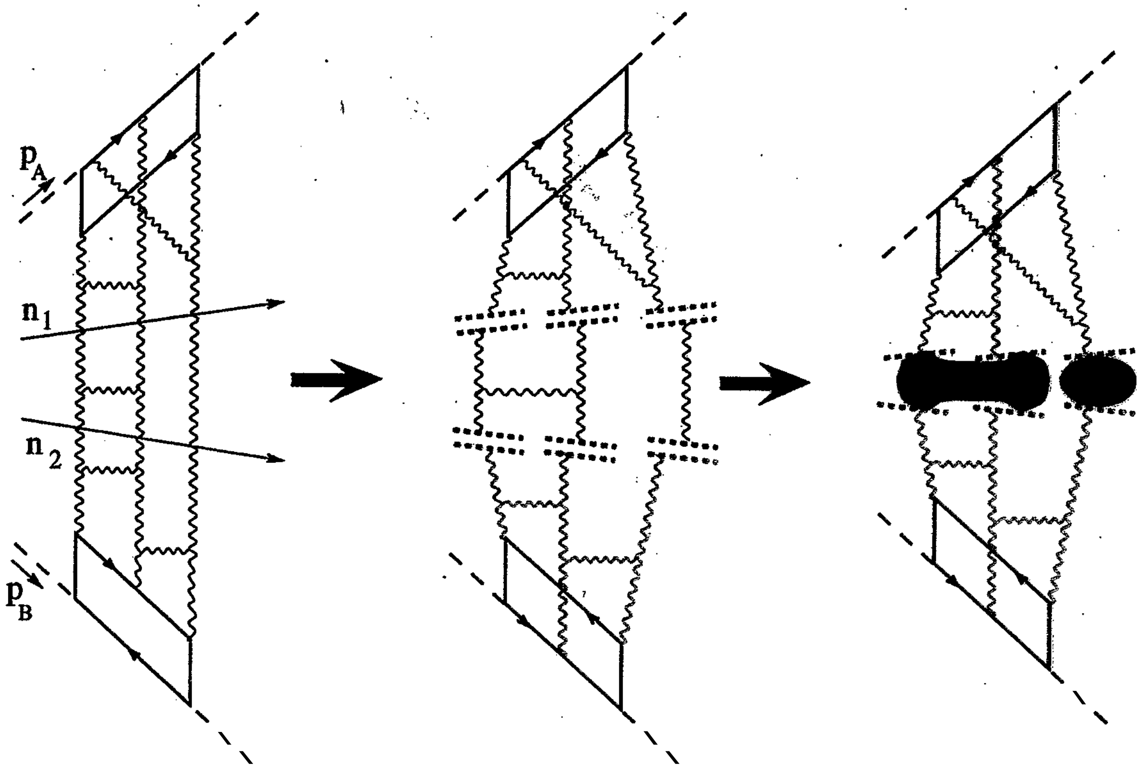
$$V(x_\perp) = [-\infty n_2 + x_\perp, \infty n_2 + x_\perp]$$

$$\Rightarrow \qquad S_{\text{eff}}(A, B) = S_{\text{eff}}(U, V)$$

$$e^{iS_{\text{eff}}(U,V)} = \int DC e^{iS(C)} e^{iS_{\text{int}}(U,C) + iS_{\text{int}}(C,V)}$$

Two factorization formulas (for rapidities η_1 and η_2) $\Rightarrow$

$$\int DAe^{iS(A)} = \int DAe^{iS(A)} \int dBe^{iS(B)}$$
$$\int dCe^{iS(C)} e^{i\int d^2x(U_i(x)U_i(x)+V_i(x)V_i(x))}$$



$\Rightarrow$ the effective action:

$$\int DAe^{iS(A)} = \int DAe^{iS(A)} \int dBe^{iS(B)} e^{iS_{\text{eff}}(U,V)}$$
$$e^{iS_{\text{eff}}(U,V)} = \int dCe^{iS(C)} e^{i\int d^2x(U_i(x)U_i(x)+V_i(x)V_i(x))}$$

In the LLA $n_1, n_2 \simeq$ light-like vectors $\Rightarrow$

$$e^{iS_{\text{eff}}(U,Y)} = \int DAe^{iS(A)}$$

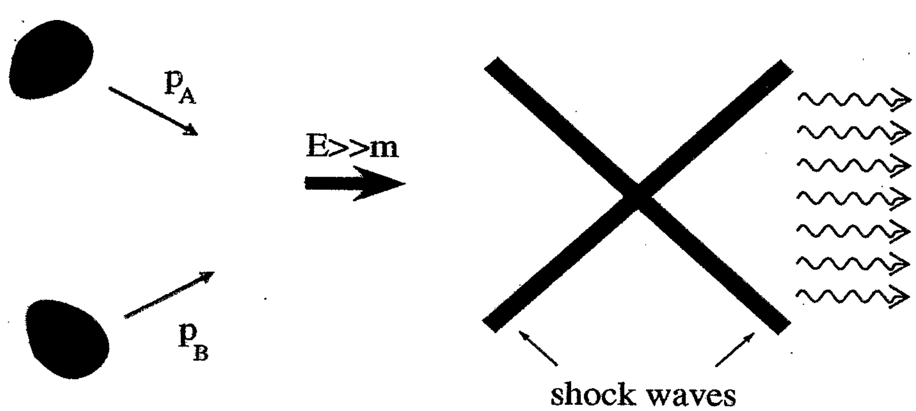
$$\exp\Big(i\int d^2xU_i(x)[-\infty n_1, \infty n_1]_x\partial_i[\infty n_1, -\infty n_1]_x$$

$$+ [-\infty n_2, \infty n_2]_x\partial_i[\infty n_2, -\infty n_2]_x V_i(x)\Big)$$

- functional integral with *two* shock-wave-type sources, U_i and V_i

Semiclassical calculation of S_{eff}: scattering of two shock waves

$$\frac{\delta}{\delta A}\Big(S(A) + \int d^2xU_iV_i(A) + W_i(A)Y_i\Big)\Big|_{A=\bar{A}} = 0$$



Approximate solution: $A = A^{(0)} + \bar{A}^{(1)} + \ldots$

$$\bar{A}_i^{(0)}(x) = U_i\theta(x_+) + V_i\theta(x_-)$$

$$\bar{A}_i^{(1)}(x) = g\int dz \frac{(x-z)_k}{(x-z)_\perp^2} \ln\left(1 - \frac{(x-z)_\perp}{x_+x_-}\right) L_{ik}(z_\perp)$$

$$\bar{A}_\pm^{(1)}(x) = \frac{g}{x_\mp}\int dz \ln\left(1 - \frac{(x-z)_\perp}{x_+x_-}\right) L(z_\perp)$$

$L = [U_i, V_i]$, $L_{ik} = [U_i, V_k]$ – Lipatov's vertex.

Effective action:

$$S_{\text{eff}}(U,V;\Delta\eta) = \int dx_\perp \left\{U_iV_i + K(U,V)g^2\Delta\eta\right\} + O(\Delta\eta^2)$$

$K(U,V) =$

$U_i(\ln\partial_\perp^2)V_i +$ $\leftarrow$ **gluon** reggeization

$L\frac{1}{\partial_\perp^2}L + L_{ik}\frac{1}{\partial_\perp^2}L_{ik} +$ $\leftarrow$ **BFKL** kernel

$L\frac{\partial_i}{\partial_\perp^2}U^\dagger\frac{\partial_k}{\partial_\perp^2}UL_{ik} + (U \leftrightarrow V)$ $\leftarrow$ 3 – pomeron vertex

Conclusions

- Factorization formula $\Rightarrow$ rigorous definition of S_{eff} fur a given $\Delta\eta$
- Semiclassical approach to S_{eff} $\Leftrightarrow$ scattering of two shock waves in QCD
- Wilson-line functional integral for S_{eff} effectively summarizes all the LLA information about high-energy scattering.

Outlook.

- Heavy-ion collisions in McLerran's model.
- Numerical calculation of the Wilson-line functional integral.
- S_{eff} in the NLO BFKL.

Transition from naive parton model to parton saturation

Jianwei Qiu

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At leading power of large momentum exchange, perturbative QCD has been very successful in interpreting and predicting high-energy scattering processes. Much of the predictive power of perturbative QCD is contained in factorization theorems [1], which provide ways to separate long- from short-distance effects in hadronic amplitude. They express nonperturbative long-distance effects in terms of universal matrix elements, which allow them to be measured experimentally or by numerical calculations. They supply systematic ways to calculate the short-distance effects perturbatively. Predictions follow when processes with different short-distance scatterings but the same nonperturbative matrix elements are compared.

On the other hand, it has been argued that for physical processes where the effective x is very small and the typical momentum exchange of the collision Q is not large the number of soft partons in a nucleus may saturate [2,3] Qualitatively, the unknown boundary of this novel regime in (z,Q) is where the conventional perturbative QCD factorization approach should fail [4].

To quantitatively identify the boundary, we try to approach it from the perturbative side by improving perturbative QCD calculations with resummed dynamical power corrections. We calculate and resum, in terms of perturbative QCD factorization approach, nuclear size enhanced power corrections to the structure functions measured in lepton-nucleus deeply inelastic scattering [5], to the centrality and rapidity dependence of single and double inclusive hadron production in proton-nucleus collisions, and to the evolution of nuclear parton distribution functions. We show that power corrections to all these quantities are expressed in terms of one universal matrix element $\langle F^{+\alpha} F_{\alpha}{}^{+} \rangle$. Our results for the Bjorken z-, Q^2- and A-dependence of nuclear shadowing in structure functions are consistent with all existing data [5]. We are in a position to predict the leading nuclear modification to nuclear parton distribution functions.

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[2] A. H. Mueller, Nucl. Phys. B **558**, 285 (1999) [arXiv:hep-ph/9904404].

[3] E. Iancu, A. Leonidov and L. McLerran, arXiv:hep-ph/0202270, and references therein.

[4] J. W. Qiu, Nucl. Phys. A **715**, 309 (2003), and references therein.

[5] J. W. Qiu and I. Vitev, Phys. Rev. Lett. **93**, 262301 (2004). [arXiv:hep-ph/0309094].

Classical and Quantum Aspects of the Color Glass Condensate
RIKEN/BNL Research Center, March 7-11,2005

Transition from naive parton model to parton saturation

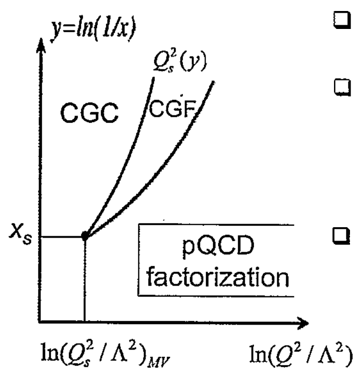
Jianwei Qiu
Iowa State University

based work done with X. Guo, M. Luo, Z. Kang, G. Stermen, I. Vitev, X. Zhang, et al.

Outline of the Talk

- ❑ Naive parton model
- ❑ **QCD** improved parton model
- ❑ Small x and coherent multiparton interactions
- ❑ Resummed power corrections to cross sections
- ❑ Resummed power corrections to DGLAP equation
- ❑ Summary and outlook

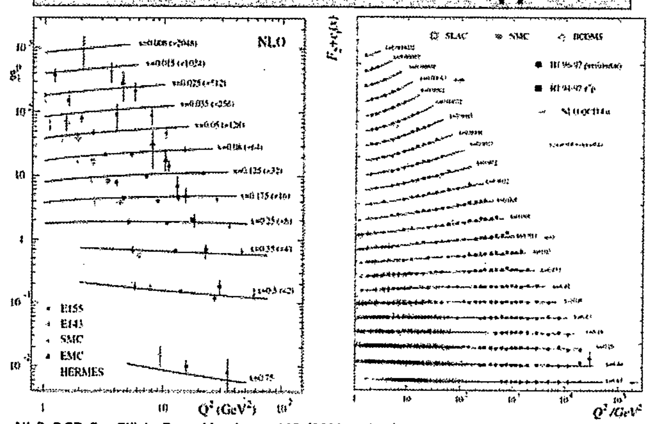
Phase diagram of parton densities



- ❑ Experiments measure cross sections, not PDFs
- ❑ PDFs are extracted based on
  - ➢ factorization
  - ➢ truncation of perturbative expansion
- ❑ How to probe the boundary between different regions?

Look for where pQCD factorization approach fails

Success of QCD factorization approach

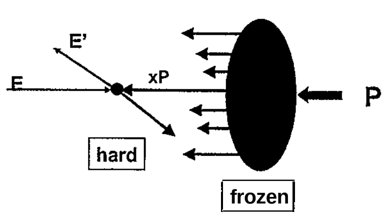


NLO QCD fit: Glück, Reya, Vogelsang, MS (2000 update)

Naïve parton model

Hard probe – Impulse approximation – Parton model

Use DIS as an example



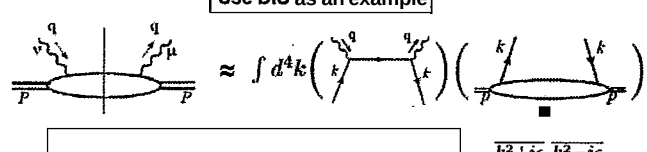
$$\sigma_{lP}(Q) \approx \int dx\, f_{q/P}(x)\, \hat{\sigma}(x,Q)$$

Convolution of two probability functions

QCD derivation of the parton model

Feynman diagrams – Pinch singularities – Factorization

Use DIS as an example



➡ a long lived parton state

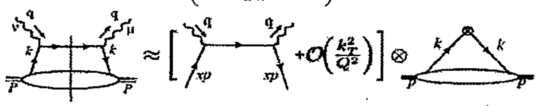
$$W^{\mu\nu}(Q^2) \approx \int \frac{dk^+}{2k^+} d^2k_T\, \hat{w}^{\mu\nu}(Q^2, k^2=0) \int dk^2 T(k) + \left(\frac{\langle k^2 \rangle}{Q^2}\right)$$

Convolution of *two* probability functions
Quantum interference between two are power suppressed

QCD improved parton model

Collinear approximation – Universal PDF's – Evolution

$$\left(k^+, k^-, k_T\right) \Rightarrow \left(k^+, \frac{k_T^2}{2k^+}, k_T\right) \Rightarrow \left(k^+, 0, 0\right)$$

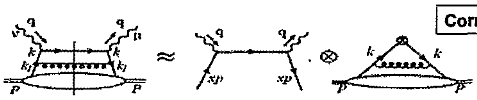


$$W^{\mu\nu}\left(Q^2\right) \approx \int dx\, \hat{w}^{\mu\nu}\left(Q^2, k = xP\right) \int d^4k\, \delta\left(x - \frac{k^+}{P^+}\right) T(k) + \left(\frac{\langle k^2\rangle}{Q^2}, \frac{\langle k_T^2\rangle}{Q^2}\right)$$

PDF's

Pinched poles in high order diagrams - evolution

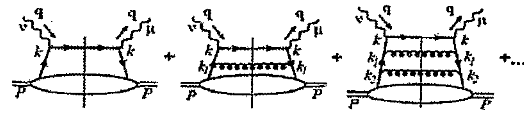
Corrections to PDF's



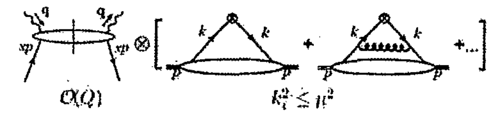
March 10, 2005 7 Jianwei Qiu, ISU

Evolution of PDFs

- perturbative pinch singularities in $\int d^4k_i$ of ladder diagrams:



- resummation of leading logarithmic contributions:



- μ^2-dependence of parton distributions ⇔ DGLAP equation

$$\mu^2 \frac{\partial}{\partial \mu^2} \phi_{i/h}(x, \mu^2) = \sum_j P_{i/j}\left(\frac{x}{x'}, \alpha_s\right) \otimes \phi_{j/h}(x', \mu^2)$$

⇒ PQCD predicts $\phi_{f/h}(x, \mu^2)$, if knows PDF's at μ_0^2

March 10, 2005 8 Jianwei Qiu, ISU

Small x and size of the hard probes

- ❑ Size of a hard probe is very localized and much smaller than a typical hadron at rest

$$1/Q \ll 2R \sim \text{fm}$$

- ❑ But, it might be larger than a Lorentz contracted hadron:

$$1/Q > 2R\left(m/p\right)$$

- ❑ low x: uncertainty in locating the parton is much larger than the size of the boosted hadron (a nucleon)

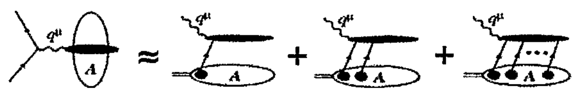
$$\frac{1}{Q} \frac{1}{xp} \gg 2R\frac{m}{P} \Rightarrow x \ll x_c \equiv \frac{1}{2mR} \approx 0.1$$

If the active *x* is small enough
a hard probe could cover several nucleons
in a Lorentz contracted large nucleus!

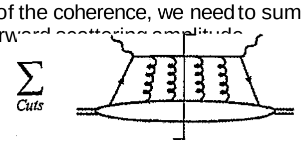
March 10, 2005 9 Jianwei Qiu, ISU

Coherent multiparton interactions

At small x, the hard probe covers several nucleons, coherent multiple scattering could be equally important at relatively low Q



To take care of the coherence, we need to sum over all cuts for a given forward scattering amplitude

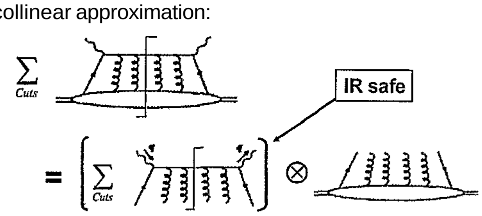


Summing over all cuts is also necessary for IR cancellation

March 10, 2005 10 Jianwei Qiu, ISU

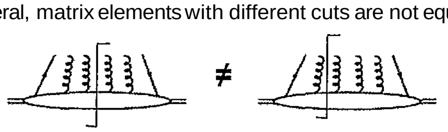
Collinear approximation is important

With collinear approximation:



IR safe

In general, matrix elements with different cuts are not equal:



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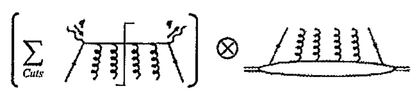
Leading contribution in medium length

Parton momentum convolution:



$$\propto \int \prod_i dy_i^-\, e^{ix_i p^+ y_i^-} \langle P_A | \prod_i F^{+\perp}\left(y_i^-\right) | P_A \rangle$$

All coordinate space integrals are localized if x is large

Leading pole approximation for dx_i integrals :

- O dx_i integrals are fixed by the poles (no pinched poles)
- ❑ x_i=0 removes the exponentials
- ❑ dy integrals can be extended to the size of nuclear matter

Leading pole leads to highest powers in medium length,
a much small number of diagrams to worry about

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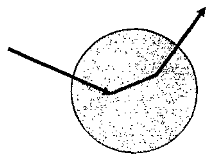
Dynamical power corrections

- Coherent multiple scattering leads to **dynamical** power corrections:



$$\frac{d\sigma^{(D)}}{d\sigma^{(S)}} \sim \alpha_s \frac{1/Q^2}{R^2} \langle F^{+\alpha} F_\alpha^{+} \rangle A^{1/3}$$

$$d\sigma \approx d\sigma^{(S)} + d\sigma^{(D)} + \ldots$$

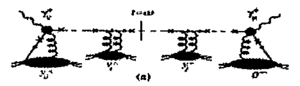
- Characteristic scale for the power corrections: $\langle F^{+\alpha} F_\alpha^{+} \rangle$
- For a hard probe: $\frac{a_s}{Q^2 R^2} \ll 1$
- **To** extract the universal matrix element, we need new observables more sensitive to $\langle F^{+\alpha} F_\alpha^{+} \rangle$

Resummation of multiple scattering

- LO contribution to DIS cross section:  $\Rightarrow \delta(x - x_B)$
- NLO contribution: 

$$\Rightarrow \underbrace{\frac{g^2}{Q^2}\left(\frac{1}{2N_c}\right)\left[2\pi^2 \tilde{F}^2(0)\right]}_{\int \frac{dy_2^- dy_1^-}{(2\pi)^2}\left[F^{+\alpha}(y_2^-)F_\alpha{}^{+}(y_1^-)\right]\theta(y_2^-)} x_B \lim_{x_1\to x} \underbrace{\left[\frac{1}{x-x_1}\delta(x_1-x_B)+\frac{1}{x_1-x}\delta(x-x_B)\right]}_{x_B\left[-\frac{d}{dx}\delta(x-x_B)\right]}$$

- **Nth order contribution:**



$$\Rightarrow \left[\frac{g^2}{Q^2}\left(\frac{1}{2N_c}\right)\left[2\pi^2 \tilde{F}^2(0)\right]\right]^N \underbrace{x_B^N \lim_{x_i\to x}\sum_{m=0}^{N}\delta(x_m-x_B)\left[\prod_{i=1}^{m}\left(\frac{1}{x_{i-1}-x_m}\right)\right]\left[\prod_{j=1}^{N-m}\left(\frac{1}{x_{m+j}-x_m}\right)\right]}_{x_B^N\left[(-1)^N \frac{1}{N!}\frac{d^N}{dx^N}\delta(x-x_B)\right]}$$

Infrared safe!

Contributions to DIS structure functions

- Transverse structure function: Qiu **and** Vitev. PRL (in **press**)

$$F_T(x_B,Q^2) = \sum_{n=0}^{N}\frac{1}{n!}\left[\frac{\xi^2}{Q^2}\left(A^{1/3}-1\right)\right]^n x_B^n \frac{d^n}{dx_B^n} F_T^{(0)}(x_B,Q^2)$$

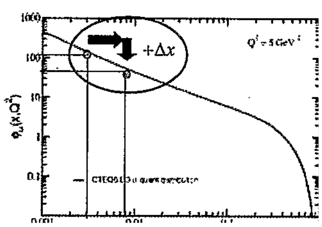
$$\approx F_T^{(0)}(x_B(1+\Delta),Q^2)$$

Similar expression of F_L

$$\Delta \equiv \frac{\xi^2}{Q^2}\left(A^{1/3}-1\right)$$

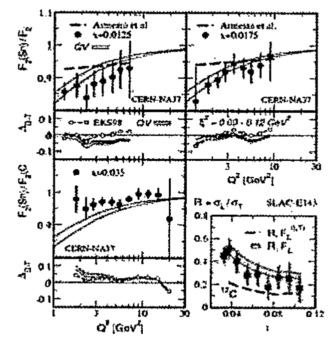
$$\xi^2 = \frac{3\pi\alpha_s}{8R^2}\langle F^{+\alpha}F_\alpha{}^{+}\rangle$$

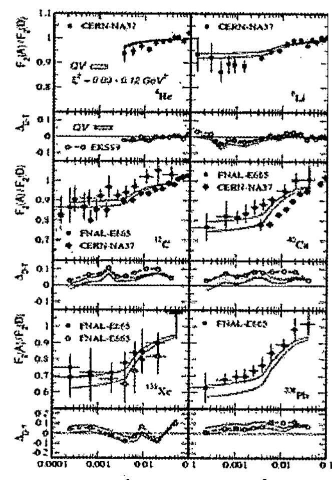
Single parameter for the power correction, and is proportional to the same characteristic scale



Neglect LT shadowing upper limit of ξ^2

$\xi^2 = 0.09 - 0.12\,\mathrm{GeV}^2$





Upper limit of $\langle F^{+\alpha} F_\alpha^{+} \rangle$ from DIS data

- Drell-Yan Q_T-broadening data:

$$\Longrightarrow \langle F^{+\alpha}F_\alpha^{+}\rangle_{DY} - 3-4 \Longrightarrow \xi^2 \approx 0.05 - 0.06\ \mathrm{GeV}^{-2}$$

- Upper limit from the shadowing data:

$$\Longrightarrow \xi^2_{Max} \approx \mathbf{0.09 - 0.12}\,\mathrm{GeV}^{-2} \Longrightarrow \langle F^{+\alpha}F_\alpha^{+}\rangle_{DIS} < 6$$

- "Saturation" scale of cold nuclear matter:

$$Q_s^2 - \xi^2 A^{1/3} \le 0.3\ \mathrm{GeV}^2 \text{ seen by quarks}$$

$$10.6\ \mathrm{GeV}^2 \text{ seen by gluons}$$

- Physical meaning of these numbers:

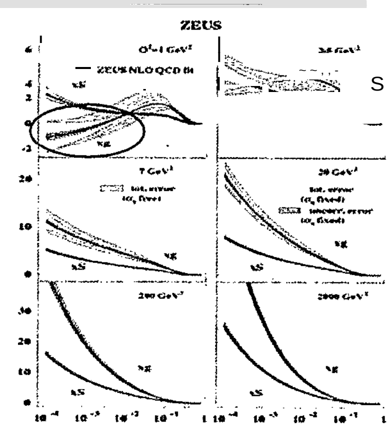
$$\langle F^{+\alpha}F_\alpha^{+}\rangle \equiv \frac{1}{p^+}\int dy_1^- \langle N|F^{+\alpha}(0)F_\alpha^{+}(y_1^-)|N\rangle\theta(y_1^-) \approx \frac{1}{2}\lim_{x\to 0} xG(x,Q^2)$$

$$\Longrightarrow \langle xG(x\to 0,Q_s^2)\rangle \le 8 \text{ in cold nuclear matter(?)}$$

Negative gluon distribution at low Q

- **NLO** global fitting based on leading twist DGLAP evolution leads to negative gluon distribution
- MRST **PDF's** have the same features

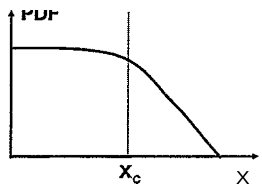
Does it mean that we have no gluon for $x < 10^{-3}$ at 1 GeV?



Leading twist shadowing

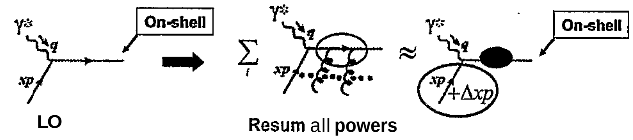
- Power corrections complement: to the leading twist shadowing:
  - Leading twist shadowing changes the x- and Q-dependence of the parton distributions
  - Power corrections to the DIS structure functions (or cross sections) are effectively equivalent to a shift in x
  - Power corrections vanish quickly as hard scale Q increases while the leading twist shadowing goes away much slower
- If leading twist shadowing is so strong that x-dependence of parton distributions saturates for $x < x_c$

additional power corrections, the shift in x, should have no effect to the cross section!

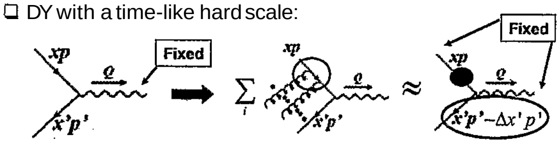


Intuition for the power corrections

- DIS with a space-like hard scale:

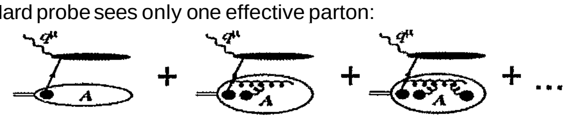


- DY with a time-like hard scale:

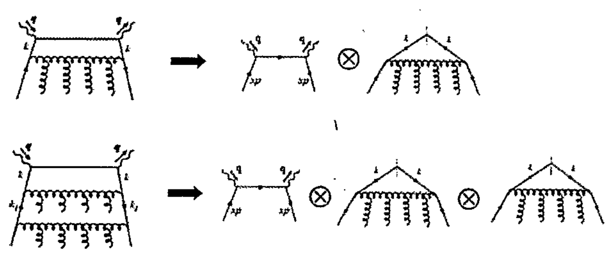


power corrections to PDFs

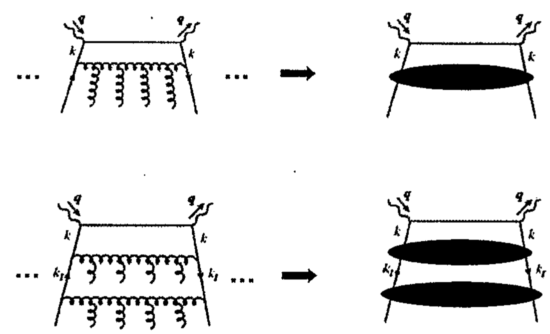
Hard probe sees only one effective parton:



Pinched poles in the ladder diagrams – corrections to evolution



Modified ladder diagrams



Modifications to DGLAP equation

Z. Kang and J. Qiu in preparation

DGLAP equation:

$$\mu^2\frac{\partial}{\partial\mu^2}\phi_{i/h}(x,\mu^2) = \sum_j P_{i/j}(\tfrac{x}{x'},\alpha_s)\otimes\phi_{j/h}(x',\mu^2)$$

What were done:

- resum all powers of leading pole coherent power corrections to all particles entering final-state
- derive a set of generalized ladder diagrams
- derive a modified DGLAP equation with the power corrections

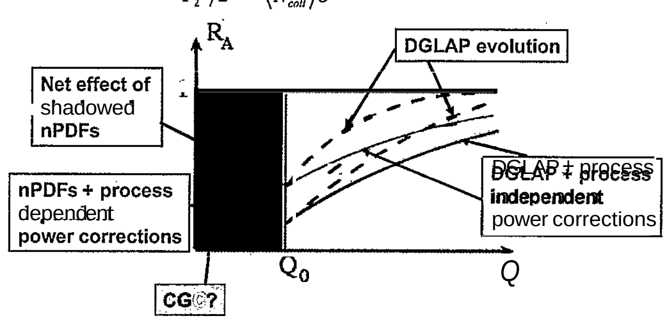
Modifications:

- shift the parton momentum fraction in PDFs in the integral part
- shift the 1/x pole by 1/(x+Δx)
- naturally generates the shadowing at low Q^2, if we evolve from high Q^2.

Role of coherent power corrections

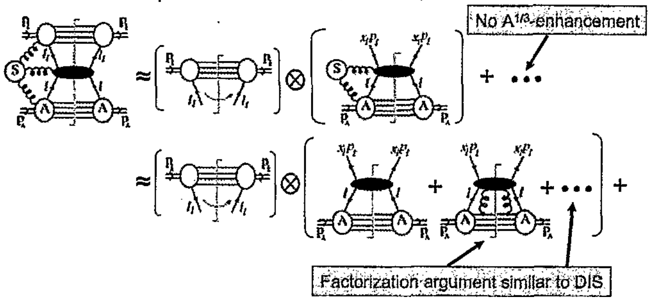
- **Ratio of physical obsewables: R_A**

$$R_A \equiv \frac{F_2^A/A}{F_2^D/2}, \quad \frac{\sigma^{dA}}{\langle N_{coll}\rangle\sigma^{NN}}, \text{ etc.}$$



Factorization in p-nucleus collisions

O A-enhanced power corrections, $A^{1/3}/Q^2$, are factorizable:



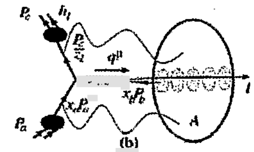
❑ But, power corrections to hard parts are process-dependent, and they are different from DIS

Power Corrections in p+A Collisions

❑ H ... fails for power corrections of the order of $/Q$ and beyond

❑ ... size enhanced dynamical power corrections in p+A could be factorized

➡ to make ...di... i... for p+A collisions



❑ **Single hadron** inclusive ...iction:

Once ...e fi... the ... g parton ... om ... e ... **and outgoing fragmentation** parton, **we uniquely fix the** momentum exchange, q^μ, and **the probe** size

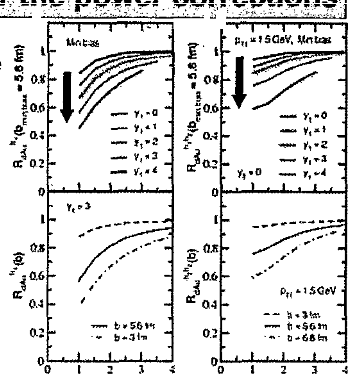
⇔ h ... along the dir... of $q^\mu - p^\mu$

Numerical results for the power corrections

❖ Similar power correction modification to single and double inclusive hadron production

- ➢ increases with rapidity
- ➢ increases with centrality
- ➢ disappears at high p_T in accord with the QCD factorization theorems
- ➢ single and double inclusive shift in $-\xi^2/t$

$$s = 2\frac{p_T^2}{z^2}(1+\cosh(y_1-y_2)),$$
$$t = -\frac{p_T^2}{z^2}(1+\exp(-y_1+y_2)),$$
$$u = -\frac{p_T^2}{z^2}(1+\exp(y_1-y_2))$$



Small at mid-rapidity C.M. energy 200 GeV
Even smaller at mid-rapidity C.M. energy 62 GeV

Qiu and Vitev, hep-ph/0405068

Our approach to multiparton interactions

❑ Advantage:

- ❖ factorization approach enables us to quantify the high order corrections
- ❖ express non-perturbative quantities in terms of matrix elements of well-defined operators – universality
- ❖ better predictive power

❑ Disadvantage:

- ❖ Rely on the factorization theorem – not easy to prove
- ❖ Hard probe might limit the region of coherence – small target

❑ Helper:

- ❖ Hard probe at small x could cover a large nuclear target

Summary and outlook

❑ Introduce a systematic factorization approach to coherent QCD multiparton interactions

❑ Leading medium size enhanced nuclear effects due to power corrections can be systematically calculated, and

❑ Identify a characteristic scale for the QCD rescattering: $\langle F^{+\alpha}F_\alpha^+\rangle$ which corresponds to a mass scale 0.6 GeV^2 (seen by gluons) in cold nuclear matter

❑ Derive coherent power corrections to DGLAP evolution equation

❑ Should be relevant for physics approaching to saturation

❑ Many applications:
jet broadening and suppression of jet correlation in p-A
...

Piu and Vitev, PLB587 (2004)
hep-ph/0405068

RIKEN BNL **Workshop,** March 7-11, 2005

Colour strings vs. Hard Pomeron

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Universidade de Santiago de Compostela, Spain

Contents:

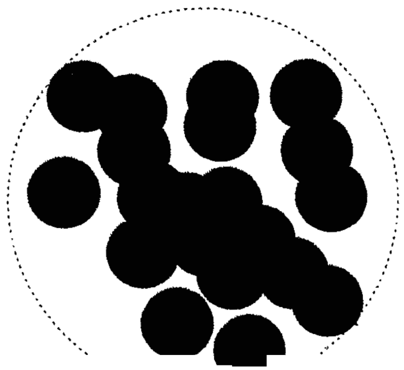
FUSING COLOUR STRINGS

- Phenomenological model for the soft region

- In a collision a certain number of colour strings are streched between the colliding partons

- Color string = strong colour field is succesively broken by creation of $q\bar{q}$ pairs

- Color strings = small areas in the transverse space filled with color field created by the colliding partons $\Rightarrow$ Phenomenon of string fusion and percolation



$$\eta = N_{st}\frac{S_1}{S_A}$$
$$S_1 = \pi r_0^2$$
$$r_0 = 0.2 \div 0.3 \text{ fm}$$
$$\eta_c = 1.1 \div 1.5.$$

- Hypothesis: clusters of overlapping strings are the sources of particle production

- For a **cluster** of n overlapping **strings** covering **an** area S_n:

Color charge of the cluster=Vectorial sum of the strings charges

$$\vec{Q}_n = \sum_{i=1}^{n} \vec{Q}_{1i}, \quad \langle \vec{Q}_{1i} \cdot \vec{Q}_{1j} \rangle = 0, \quad \vec{Q}_n^2 = n\vec{Q}_1^2, \tag{1}$$

$$Q_n = \sqrt{\frac{nS_n}{S_1}} Q_1, \quad \mu_n = \sqrt{\frac{nS_n}{S_1}} \mu_1, \quad \langle p_T^2 \rangle_n = \sqrt{\frac{nS_1}{S_n}} \langle p_T^2 \rangle_1. \tag{2}$$

For strings without interaction: $S_n = nS_1$, $Q_n = nQ_1 \rightarrow \mu_n = n\mu_1$, $\langle p_T^2 \rangle_n = \langle p_T^2 \rangle_1$

For strings with max overlapping: $S_n = S_1$, $Q_n = \sqrt{n} Q_1 \rightarrow \mu_n = \sqrt{n}\mu_1$, $\langle p_T^2 \rangle_n = \sqrt{n}\langle p_T^2 \rangle_1$

- Moreover, one **can obtain** the analytic **expression:**

$$< \frac{nS_1}{S_n} >= \frac{\eta}{1 - \exp(-\eta)} \equiv \frac{1}{F(\eta)^2} \tag{3}$$

so

$$\mu = N_{strings}\, F(\eta) \mu_1 \ , \qquad < p_T^2 >= \frac{1}{F(\eta)} < p_T^2 >_1 \tag{4}$$

PERTURBATIVE QCD POMERON

Inclusive cross section in pQCD, taking $A = B$ and constant nuclear density for $|b| < R_A$:

$$I_A(y, k) = A^{2/3}\pi R_0^2 \frac{8N_c\alpha_s}{k^2}\int d^2 r e^{ikr}[\Delta\Phi_A(Y - y, r)][\Delta\Phi_A(y, r)], \qquad (5)$$

where $\Phi(y, r)$ is the sum of all fan diagrams connecting the pomeron at rapidity y and of the transverse dimension r with the colliding nuclei, one at rest and the other at rapidity Y.

In the momentum space, function $\phi_A(y, r) = \Phi(y, r)/(2\pi r^2)$ satisfies

$$\frac{\partial\phi(y, q)}{\partial\bar{y}} = -H\phi(y, q) - \phi^2(y, q), \qquad (6)$$

where $\bar{y} = \bar{\alpha}y$, $\bar{\alpha} = \alpha_s N_c/\pi$, α_s and N_c are the strong coupling constant and the number of colours, respectively, and H is the BFKL Hamiltonian

This equation has to be solved with the initial condition at $y = 0$ determined by the colour dipole distribution in the nucleon smeared by the profile function of the nucleus.

We take the initial condition in accordance with the Golec-Biernat distribution:

$$\phi(0, \boldsymbol{q}) = -\frac{1}{2} a \, \mathrm{Ei}\left(-\frac{q^2}{0.3567\,\mathrm{GeV}^2}\right), \qquad (7)$$

with

$$a = A^{1/3} \frac{20.8\,\mathrm{mb}}{\pi R_0^2} \qquad (8)$$

Evolving $\phi(y, q)$ up to values $\bar{y} = 3$ we found the inclusive cross-section at center rapidity for energies corresponding to the overall rapidity $Y = \bar{Y}/\bar{\alpha}$. With $\bar{\ } = 6$ and $\alpha_s = 0.2$ this gives $Y \sim 31$. The overall cutoffs for integration momenta in Eq.(26) were taken according to $0.3.10^{-8}$ GeV/c $< q < 0.3.10^{+16}$ GeV/c.

At relatively small momenta the inclusive cross-sections are proportional to A, that is to *the number of participants*

At larger momenta they grow with A faster, however noticeably slowlier than the number of collisions, approximately as $A^{1.1}$

The interval of momenta for which $I_A \propto A$ is growing with energy, so that one may conjecture that at infinite energies all the spectrum will be proportional to A

New developments in the dipole model

Michael Lublinsky

University of Connecticut

E. Levin and M. L., hep-ph/0501173;
Phys. Lett. **B607** (2005) 131;
Nucl. Phys. **A730** (2004) 191;

A. Kovner and M. L., hep-ph/0502071.

M. Lublinsky

New Developments in the Dipole model.

Abstract: The dipole model of high energy QCD is reformulated as a linear functional evolution for generating functional. Using this formalism we rederive the Balitsky's hierarchy and the JIMWLK equation. We extend the dipole model by including the dipole recombination process. This allows us to introduce a new effective theory for Pomeron interactions in the large Nc and high energy limits,

$$\frac{\partial N(Y)}{\partial Y_0} = 0\,; \qquad \frac{\partial\, Z[s]}{\partial\, Y} = \chi[s]\; Z[s]$$

$$\boxed{\partial\, W^t(Y_0,\,[s])\,/\,\partial\, Y_0 \;=\; \chi[s]\;W^t(Y_0,\,[s])}$$

Jalilian-Marian - Iancu - McLerran - Weigert - Leonidov - Kovner

$$\chi[s] \;=\; \boldsymbol{-2}\,[\mathrm{s} - \mathrm{ss}]\,\frac{\mathbf{S}}{s\,s} \;-\; \boldsymbol{-2}\,\frac{\delta}{\delta\, s}\,[\mathrm{s} - \mathrm{ss}]$$

$$\boxed{\chi^{JIMWLK} \;=\; \chi[s] \;+\; \tfrac{1}{N_c^2}\,\chi^{cc}}$$

$$\chi^{cc} \;\propto\; \frac{\delta^2}{\delta\, s\,\delta\, s}$$

A Kovner and M.L., hep-ph/0502071

R. Janik, hep-ph/0409256

Functional evolution with Pomeron loops

E Levin and M.L., hep-ph/0501173

$$\frac{\partial\, Z[u]}{\partial\, Y} \;=\; \chi[u]\;Z[u]$$

$$\boxed{\chi[u] \;=\; \chi^{1\,\rightarrow\,2}[u] \;+\; \chi^{2\,\rightarrow\,1}[u]}$$

$$\chi^{1\,\rightarrow\,2} \;=\; -\,V_{1\,\rightarrow\,1}[u] \;+\; V_{1\,\rightarrow\,2}[u]$$

$$\chi^{2\,\rightarrow\,1} \;=\; -\,V_{2\,\rightarrow\,2}[u] \;+\; V_{2\,\rightarrow\,1}[u]$$

$$V_{1\rightarrow 2}[u] \;=\; \int dPS\,\Gamma(1\rightarrow 2)\;\;u\,u\,\frac{\delta}{\delta u}\,.$$

$$V_{1\,\rightarrow\,1}[u] \;=\; \int\; dPS\,\Gamma(1\rightarrow 2)\;\;u\,\frac{\delta}{\delta\, u}.$$

$$V_{2\,\rightarrow\,1}[u] \;=\; \int\; dPS\;\Gamma(2\rightarrow\;\mathbf{1})\,u\,\frac{6}{\delta\,u}\,\frac{6}{\delta\,u}\,.$$

$$V_{2\,\rightarrow\,2}[u] \;=\; J\;\; dPS\;\Gamma(2\rightarrow\;\mathbf{1})\,u\,u\,\frac{\delta}{\delta\,u}\,\frac{\delta}{\mathbf{su}}\,.$$

M. Lublinsky

$$\frac{\partial}{\bar{\alpha}_s\,\partial\,Y} \qquad \hat{K}^T \quad \gamma_n \; - \; \gamma_{n+1} \,] + \hat{\kappa}^{\cdot} \; \gamma_{n-1}$$

$$\frac{\partial\,\rho_n^p}{\bar{\alpha}_s\,\partial\,Y} \; = \; \hat{K} \; \left[\,\rho_n^p \; + \; \rho_{n-1}^p\,\right] - \hat{\kappa}\,\rho_{n+1}^p$$

$$\frac{\partial\,\rho_r^t}{\bar{\alpha}_s\,\partial\,Y} \qquad K \; \left|\,\rho_n \; + \; \rho_{n-1}\,\right| - \kappa\,\rho_{n+1}$$

$$\boxed{\Gamma\,(2 \rightarrow 1)}$$

$$\frac{\partial\,\rho_1^t}{\partial\,Y} \;\sim\; \Gamma\,(2 \rightarrow 1)\;\rho_2^t$$

$$\frac{\partial\,\gamma_1}{\bar{\alpha}_s\,\partial\,Y} \;=\; \int \Gamma(1 \rightarrow 2)\;[\,\gamma_1 \;-\; \gamma_2\,]$$

$$\gamma_1\,(x,y) \;=\; \int \sigma_{BA}(x,y;\,\bar{x},\bar{y})\;\rho_1^t(\bar{x},\bar{y})\,d^2\bar{x}\,d^2\bar{y}$$

$$\gamma_2\,(x_1,y_1;\,x_2,y_2) \;=\; \int \sigma_{BA}(x_1,y_1;\,\bar{x}_1,\bar{y}_1)\,\sigma_{BA}(x_2,y_2;\,\bar{x}_2,\bar{y}_2)$$

$$\times\;\rho_2^t(\bar{x}_1,\bar{y}_1;\,\bar{x}_2,\bar{y}_2)\,d^2\bar{x}_1\,d^2\bar{y}_1\,d^2\bar{x}_2\,d^2\,\bar{y}_2$$

$$\Gamma_{2\rightarrow 1}\,(1+2 \rightarrow x,y) \;=$$

$$\frac{2N_c^2}{\alpha_s^2}\;\wedge\quad\wedge\quad \int_{1'\,2'} \Gamma_{1\rightarrow 2}\;(\mathrm{x},\mathrm{y} \rightarrow 1'+2')\;\;\sigma_{BA}\left(1;1'\right)\;\;\sigma_{BA}\left(2;2'\right)$$

$\Delta_x\,\Delta_y\;\int_{1'\,2'}$ can be worked out, down to the computer ready *ex*pression

M. Lublinsky

Summary

- High energy evolution in QCD can be successfully described by a classical branching process with conserved probabilities.

- The linear functional evolution for generating functional is an efficient tool capable to accommodate most of the nonlinear dynamics.

- The dipole merging process can be successfully introduced leading to an effective theory which governs Pomeron dynamics in QCD at high energy, in the leading logarithmic approximation, and in the limit where N_c, the number of colors, is large.

Can Hydrodynamic Description of Heavy Ion Collisions be Derived from Feynman Diagrams?

Yuri V. Kovchegov

Department of Physics, The Ohio State University
Columbus, OH 43210

This talk is based on the paper [1].

We consider the problem of thermalization in heavy ion collisions. Thermalization in heavy ion collisions in the weak coupling framework can be viewed as a transition from the initial state Color Glass Condensate dynamics, characterized by the energy density scaling like $\epsilon \sim 1/\tau$ with τ the proper time, to the hydrodynamics-driven expansion of the quark-gluon plasma with $\epsilon \sim 1/\tau^{4/3}$ or higher power of $1/\tau$ for the boost non-invariant case. (Of course at realistic temperatures achieved in heavy ion collisions the power of 4/3 may become somewhat smaller: however, it is always greater than 1 for hydrodynamic expansion.) In this talk we argue that, at any order of the perturbative expansion in the QCD coupling constant, the gluon field generated in an ultrarelativistic heavy ion collision leads to energy density scaling as $\epsilon \sim 1/\tau$ for late times $\tau \gg 1/Q_s$. Therefore it is likely that thermalization and hydrodynamic description of the gluon system produced in heavy ion collisions can not result from perturbative QCD diagrams at these late times.

It may be possible that corrections to the saturation/Color Glass initial conditions would contribute towards modifying the $\epsilon \sim 1/\tau$ scaling to some higher power. Thus one should be interested in Feynman diagrams which would bring in τ-dependent corrections to $\epsilon \sim 1/\tau$ scaling of the (classical) gluon fields in the initial stages of the collisions (see slide 1). Unfortunately, after examining a number of diagrams, we noticed that while many of them introduce τ-dependent corrections to the initial conditions, such corrections are subleading and small at large τ and do not modify $\epsilon \sim 1/\tau$ scaling at late times. After reaching this conclusion we have constructed a general argument proving that $\epsilon \sim 1/\tau$ scaling always dominates at late times, which we are presenting here.

We begin by considering the most general case of boost-invariant gluon production (see slide 2). We argue that $\epsilon \sim 1/\tau$ scaling persists to all orders in the coupling constant α_s (slide 3). The argument is based on a simple observation that τ-dependent corrections to the classical gluon field may only come in through powers of gluon virtuality k^2 in momentum space with each power of k^2 giving rise to a power of $1/\tau$ (slide 4). In order for the on-mass shell amplitude (at $k^2 = 0$) to be non-singular only positive powers of k^2 axe allowed: hence, the corrections come in only as inverse extra powers of τ and are negligible at late times. We proceed by generalizing our results to rapidity-dependent distributions. The $\epsilon \sim 1/\tau$ scaling does not get modified by rapidity-dependent corrections either. Rapidity-dependent corrections come in as, for example, powers of k_+, which is one of light cone components of the gluon's momentum. However, powers of k_+ do not modify the τ-dependence of energy density. We also show that $\epsilon \sim 1/\tau$ scaling persists even when massless quarks are included in the problem. Therefore it appears that perturbative thermalization can not happen in heavy ion collisions. We try to give a physical explanation in slide 5. We conclude by arguing that if perturbative thermalization is impossible, than the non-perturbative QCD effects must be responsible for the formation of quark-gluon plasma (QGP) at RHIC. Such non-perturbative effects could be due to the infrared modes with momenta of the order of Λ_{QCD}. They can also be due to the non-equilibrium analogue of the ultra-soft modes of finite temperature QCD: those modes have momenta of the order of $g^2 T$ with T the temperature of the quark-gluon plasma. The dynamics of these modes is known to be non-perturbative and may contribute to thermalization.

[1] Yu. V. Kovchegov, hep-ph/0503038.

Our Approach

Can one find diagrams giving gluon fields which would lead to energy density scaling as

$$\varepsilon \sim \frac{1}{\tau^{1+\Delta}}, \qquad \Delta > 0 \qquad ?$$

Classical fields give energy density scaling as

$$\varepsilon^{classical} \sim \frac{1}{\tau}$$

Can quantum corrections to classical fields modify the power of tau (in the leading late-times asymptotics)? Is there analogues of leading log resummations (e.g. ln t), anomalous dimensions?

Energy-Momentum Tensor of a General Gluon Field

Let us start with the most general form of the gluon field

$$A_\mu^a(x) = \int \frac{d^4k}{(2\pi)^4} e^{-ik\cdot x} \frac{-i}{k^2 + i\varepsilon k_0} J_\mu^a(k)$$

plug it into the expression for the energy momentum tensor

$$T^{\mu\nu} = \left\langle - F^{a\,\mu\rho} F^{a\,\mu}{}_{\rho} + \frac{1}{4} g^{\mu\nu} (F^{a}_{\rho\sigma})^2 \right\rangle$$

keeping only the Abelian part of the energy-momentum tensor for now.

Energy-Momentum Tensor of a General Gluon Field

When the dust settles we get

$$\varepsilon = \frac{\pi}{2}\int d^2k \frac{dN}{d^2k d^2b d\eta} k_T^2 \{[J_1(k_T\tau)]^2 + [J_0(k_T\tau)]^2\}$$

leading to

$$\varepsilon(\tau >> 1/Q_S) \approx \frac{1}{\tau}\int d^2k \frac{dN}{d^2k d^2b d\eta} k_T = \frac{1}{\tau}\frac{dE_T}{d^2b d\eta}$$

We have established that e has a non-zero term scaling as 1/t .

But how do we know that it does not get cancelled by the rest of the expression, which we neglected by putting $k^2=k'^2=0$ in the argument off, ?

Energy-Momentum Tensor of a General Gluon Field

For a wide class of amplitudes we can write

$$f_1(k^2, k'^2, k_T) - f_1(k^2 = 0, k'^2 = 0, k_T) = (k^2 k'^2)^\Delta g(k^2, k'^2, k_T)$$

with $\quad g(k^2 = 0, k'^2 = 0, k_T) \neq 0$

Then, using the following integral:

$$\int_{-\infty}^{\infty} dk_+ \, dk_- \, e^{-ik_+ x_- - ik_- x_+} (k^2 + i\varepsilon k_0)^{\Delta-1} = -\frac{2\pi^2}{\Gamma(1-\Delta)} \left(\frac{2k_T}{\tau} \right)^\Delta e^{i\pi\Delta} J_{-\Delta}(k_T \tau)$$

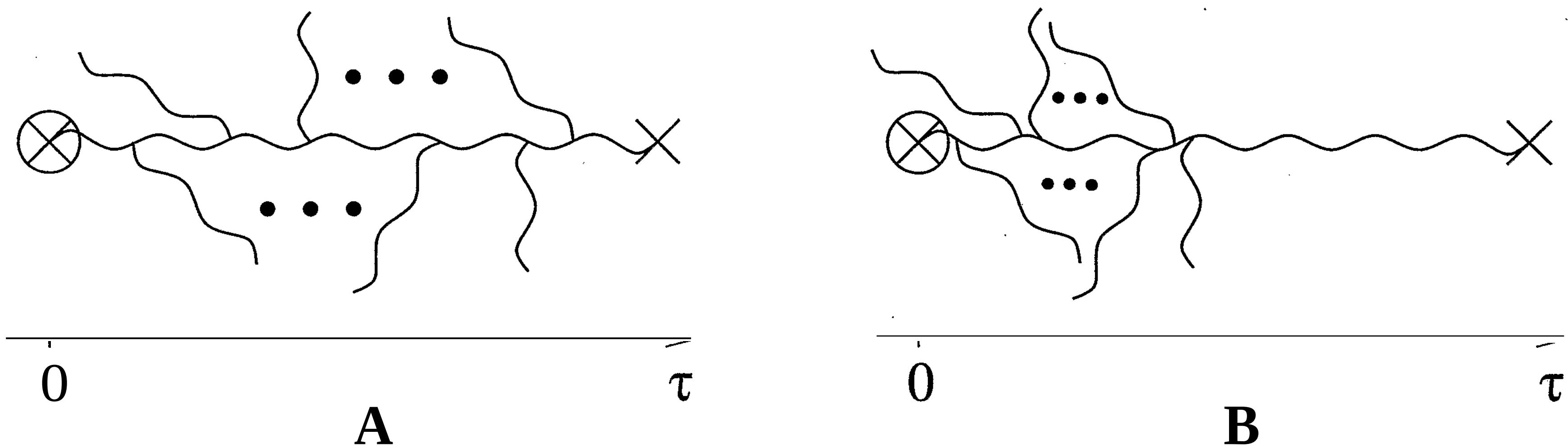
we see that each positive power of k^2 leads to a power of 1/t , such that the neglected terms above scale as

$$\sim \frac{1}{\tau^{1+2\Delta}}, \qquad \Delta > 0$$

They are subleading at large t and do not cancel the leading 1/t term.

Physical Interpretation

Is it "free streaming"?



A general gluon production diagram. The gluon is produced and multiply rescatters at all proper times.

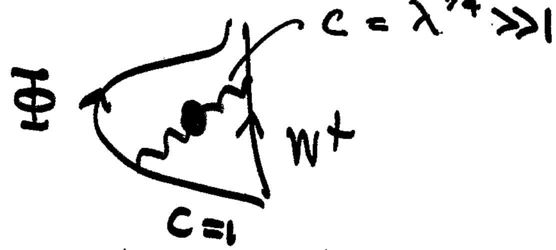
The dominant contribution appears to come from all interactions happening early.

→ Not free streaming in general, but free streaming dominates at late times.

SQGP = Reality from Duality

Jiz / BNL 05'

● Bound states



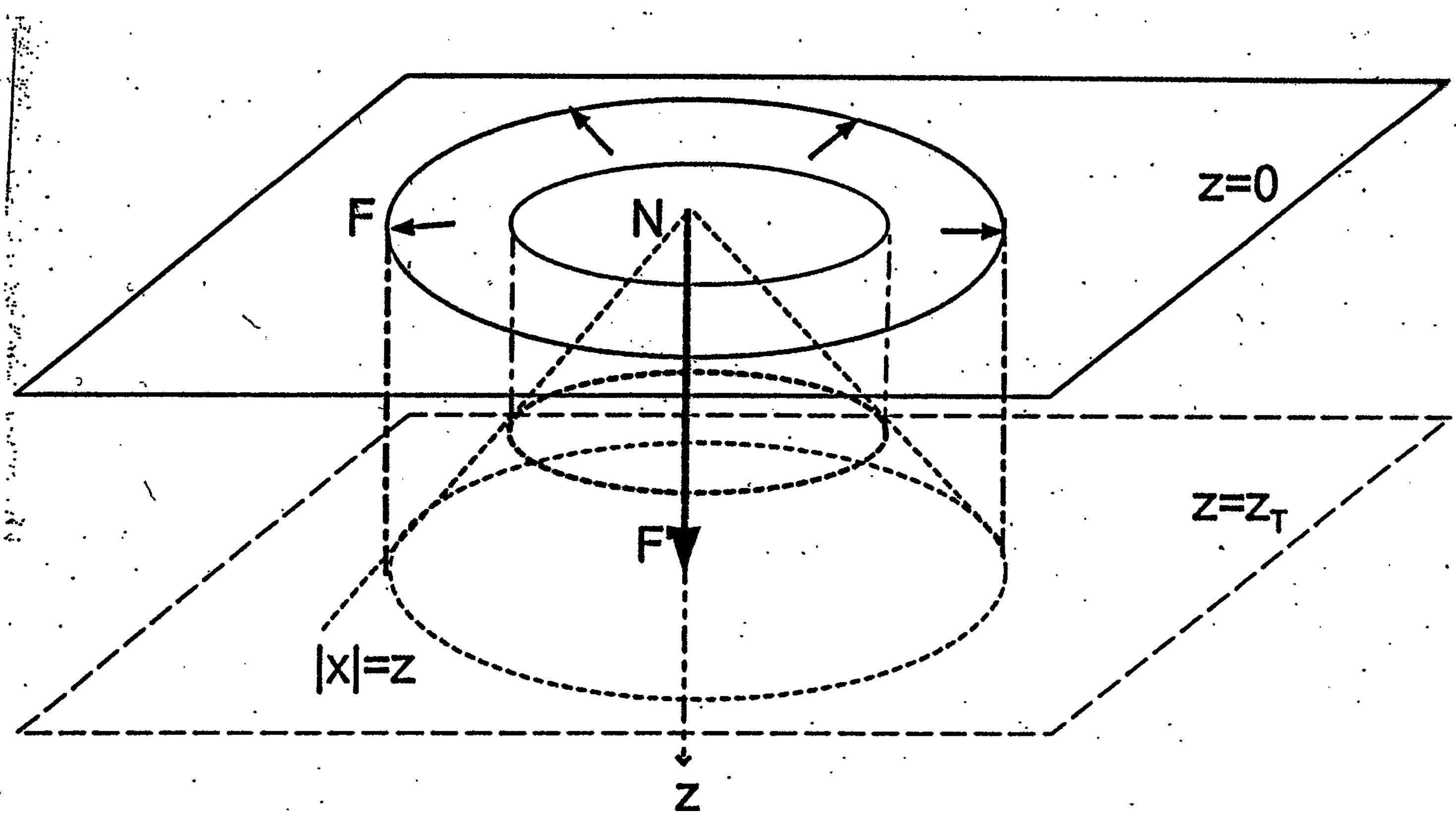
(ΦW^+) potential like ! $(-\sqrt{\lambda}/r)$

$$E_{nl} = \pm \delta m \left(1 + \left(\frac{c}{n + 1/2 + \sqrt{l^2 - c^2}} \right)^2 \right)^{-1/2}$$

$$c = \sqrt{\lambda} \gg 1 \quad : \quad E_{nl} \simeq \pm \underbrace{\frac{\delta m}{c}}_{\pm T} \left((n + 1/2) + \underbrace{(l^2 - c^2)^{1/2}}_{c_\Delta} \right)$$

Shuryak + Zahed 03'

Jet Quenching: Color opacity



$$ds^2 = \frac{\alpha' R^2}{z^2}\left(-\left(1-\frac{z^4}{z_T^4}\right)dt^2 + d\vec{x}^2 + \frac{dz^2}{(1-z^4/z_T^4)}\right)$$

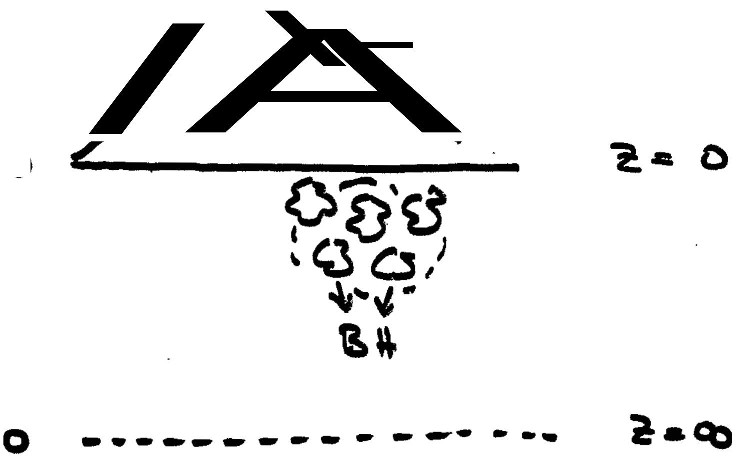
in bulk: $$dt^2 = \frac{dz^2}{(1-z^4/z_T^4)^2}$$

on Bound: $$dt^2 = \frac{d\vec{x}^2}{(1-z^4/z_T^4)}$$

$$\frac{d\vec{x}}{dt} = 1 - \frac{x_\perp^4}{z_T^4} = 1-(\pi T x)^4$$ Stalls!

Sin + Zahed 04'

Black-Hole dual of RHIC Fireball



$$ds^2 = -f\,dt^2 + f^{-1}dr^2 + d\vec{x}^{\,2} + r^2 d\Omega_5$$

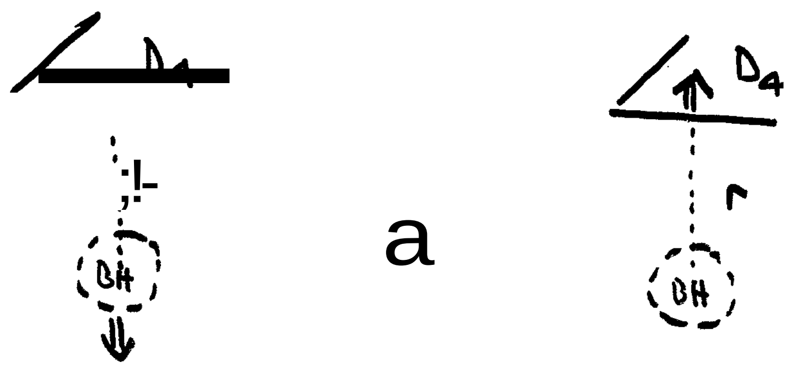
$$f = 1 - \frac{G_5 M}{r^2} + \left(\frac{r}{R}\right)^2$$

$$r_{BH} = (R^2 G_5 M)^{1/4}$$

$$Q_S^{-1} \lesssim r_{BH}$$

Shuryak + Sin + Zahed 05'

Cooling from AdS



induced metric on D_4 is Robertson-Walker like:

$$ds^2 = -d\tau^2 + a^2 d\vec{x}^2$$

$$\left(\frac{\dot{a}}{a}\right)^2 = \frac{a}{16}\,\frac{1}{a^{2/3}}\left(\left(\frac{E}{a^4}+1\right)^2 - \left(1 - \frac{r_{BH}^3}{R^3}\frac{1}{a^4}\right)\right)$$

- Cooling Contd.

$$a \uparrow \; : \quad \dot{a} \simeq a^{-4/3}$$

$$a(\tau) \simeq \tau^{\frac{3}{7}}$$

$$T(\tau) = \frac{T_{BH}}{\tau^{.43}}$$

Cool faster than Bjorken : $\frac{T_i}{\tau^{.33}}$

Shuryak + Sin + Zahed 05'

Observational constraints on Q_s from cosmic ray airshower data (hep-ph/0408073, 0501165)

A. Dumitru (ITP, Frankfurt U.)

- Cosmic Ray Airshowers are h+A collisions
- Primary Energy up to $\sim 10^{11}$ GeV ($\sqrt{s} \simeq 350$ TeV)
- Xmax mainly sensitive to forw. region and "small" tr. mom.
- BUT: small target nuclei (^{14}N), min. bias

Summary:

- **Cosmic ray airshowers are sensitive to QCD evolution scenario.**
- **Indications for a less rapid growth of $Q_s(x)$ as compared to RHIC or HERA.**
- **High-density effects increase inelasticity (forward suppression) --> hadron-induced showers resemble those of "nuclei" in present models (X_{max} lower, more μ, ...)**
- **Lighter Composition predicted**

running coupl BFKL : $\bar{\alpha}_s(Q^2) = b_0 / \log(Q^2/\Lambda^2)$

$$Q_s^2 = \Lambda^2 \exp\left(\log(Q_0^2/\Lambda^2)\sqrt{1+2c\,\bar{\alpha}_s\,y}\right)$$

yP	10.7	17.3	26.1
Qs r.c.	1.1 GeV	2.4 GeV	5.9 GeV
Qs f.c.	1.4 GeV	4.5 GeV	19.2 GeV

$\lambda = 0.28$; **central** "p+N"; (Q_0/Λ)^2~Nval/3

OHH

Quark-Nucleus Scattering

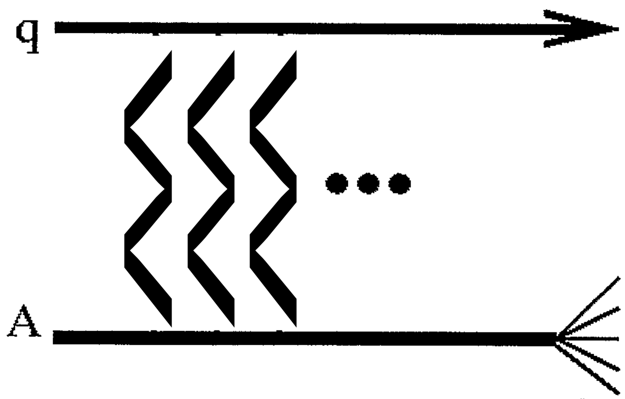
Quark Scattering Amplitude : $\langle q , p \rangle = \bar{u}(q)\,\tau(q,p)\,u(p)$

with $\tau(q,p) = 2\pi\,\delta(p-q)\,\gamma \int d^2x_t \left[V(x_t) - 1 \right] e^{i x_t (q_t - p_t)}$

$$V(x_t) = P \exp\left(-i g^2 \int_{\infty}^{\infty} dx \frac{1}{\partial_t^2} \rho^a(x, x_t)\, t^a\right)$$

Color averaging with a Gaussian (MV) leads to ($q_t > 0$):

$$C(q_t) = \int \frac{d^2 r_t}{(2\pi)^2} e^{i q_t r_t} \exp\left[-2 Q_s^2 \int_{\Lambda} \frac{d^2 p_t}{(2\pi)^2} \frac{1}{p_t^4} \left(1 - e^{i p_t r_t}\right) \right]$$



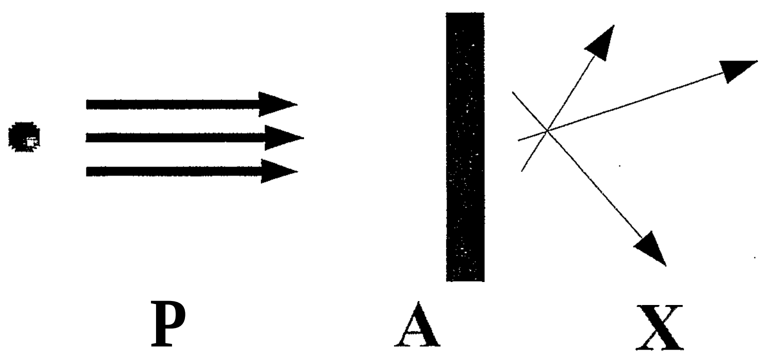
AD + J. Jal.-Marian: PRL 89 (2002

Shattering the proton

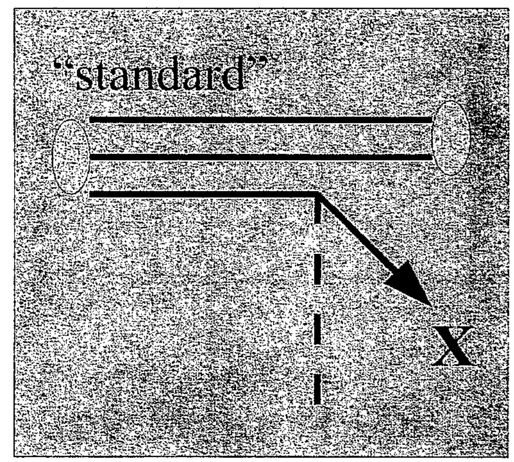
Probability for quark to be scattered to $q_t \sim 0$ (<u>with</u> color exchange !) :

$$\int_0^{\Lambda} d^2 q_t \frac{d\sigma^{inel}}{d^2 b \; d^2 q_t} \simeq \frac{\pi \Lambda^2}{Q_s^2 \log\left(Q_s / \Lambda\right)}$$

--> suppression of "beam-jet remnants'' (soft physics) in the BBL



$$\langle q_t \rangle \sim Q_s$$



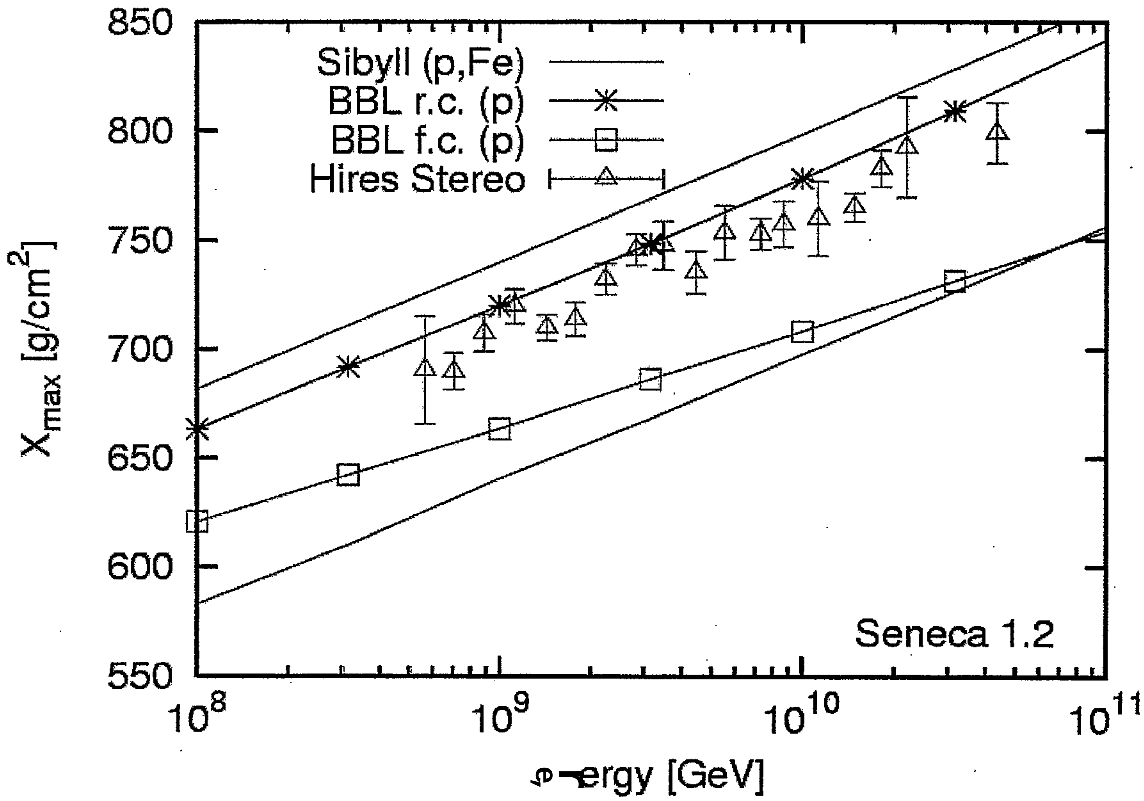
<u>**All partons resolved at scale Qs, coherence of proton destroyed completely.**</u>

Gerland, Strikman, AD: PRL 90 (2003)

Cosmic Ray Airshowers

hep-ph/0408073

"*penetration depth*"



=> Sensitive to evolution of Qs !!

⭑ $1/x^{0.3}$ yields too large Qs at GZK energies

Saturation, Unitarity and Fluctuations in High Energy Collisions

A. H. Mueller

Abstract of Talk:

This talk covers the essential equations governing high-density QCD and its application to small-x processes as well as the early stages of high-energy heavy ion reactions.

The simplest equation which imposes unitarity is the Kovchegov, a sort of mean field equation for high-energy scattering. However, fluctuations allow unitarity-violating processes to occur which are not suppressed by the Kovchegov equation. Such difficulties are avoided by carefully considering the evolution in the low-density region.

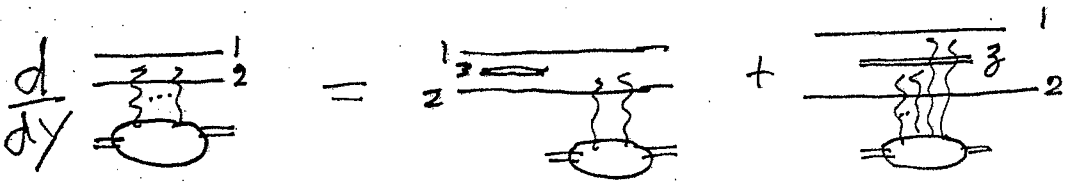
Such a careful consideration leads one to realize that there is a close connection between rapidity evolution in small-x QCD and time-evolution in reaction-diffusion equations in statistical mechanics. **A** simple statistical model, the Brunet-Derrida model, is briefly described and the correspondence to the essential elements in QCD evolution near the unitarity limit are noted.

Finally, it is emphasized that a simple description of QCD evolution near the unitarity limit is only possible if evolution is carried out in an event by event way rather than by dealing with average quantities.

Saturation, Unitarity and Flustuations in High Energy Collisions

1. Mean field approximation

(a) The Kovchegov equation



$$\frac{d}{dy} S(\underline{x}_1-\underline{x}_2, b) = \frac{\alpha N_c}{2\pi^2}\int \frac{d^2z\,(\underline{x}_1-\underline{x}_2)^2}{(\underline{x}_1-\underline{z})^2(\underline{x}_2-\underline{z})^2}\left[S^{(2)}(\underline{x}_1-\underline{z}, \underline{x}_2-\underline{z}, b) - S(\underline{x}_1-\underline{x}_2, \underline{b})\right]$$

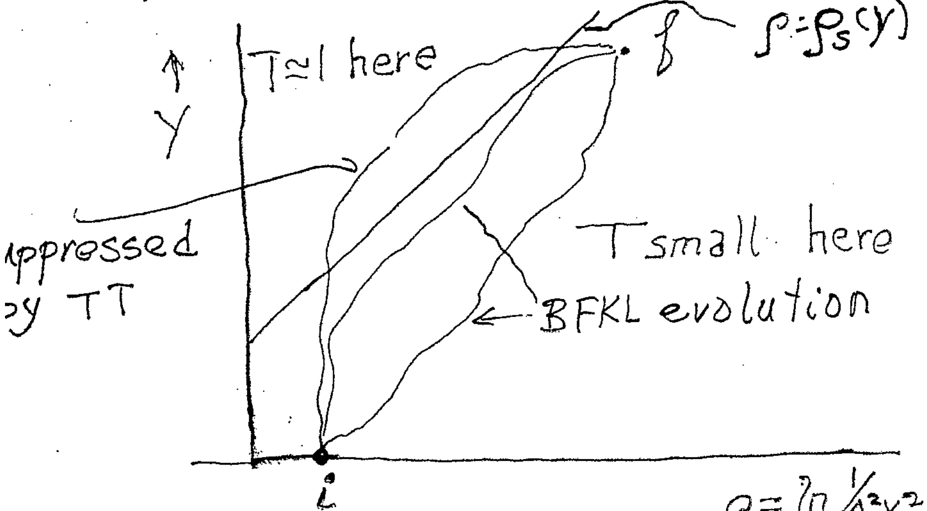
$$\downarrow$$

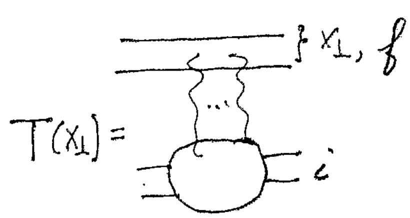
$$S(\underline{x}_1-\underline{z})\cdot S(\underline{x}_2-\underline{z})$$

$$\frac{d}{dY} T(\underline{x}_1-\underline{x}_2) = \frac{\alpha N_c}{2\pi^2}\int \frac{d^2z\,(\underline{x}_1-\underline{x}_2)^2}{(\underline{x}_1-\underline{z})^2 (\underline{x}_2-\underline{z})^2}\left[T(\underline{x}_1-\underline{z}) + T(\underline{z}-\underline{x}_2) - T(\underline{x}_1-\underline{x}_2) - T(\underline{x}_1-\underline{z})\,T(\underline{x}_2-\underline{z})\right]$$

Kovchegov equation

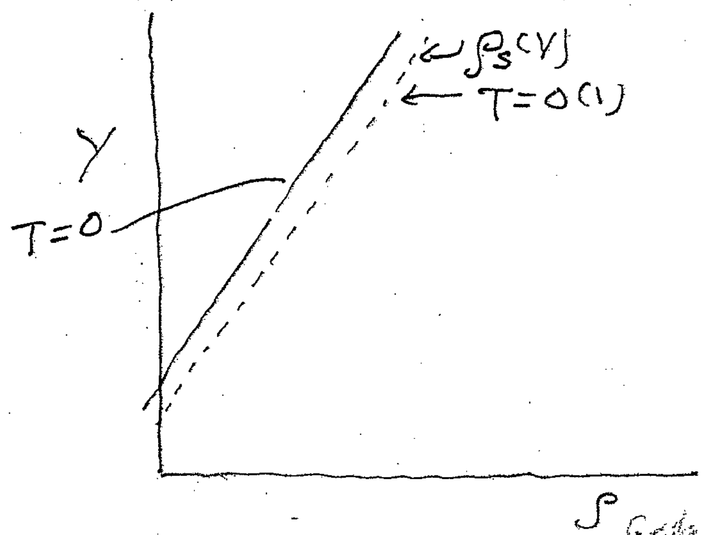
(b) The saturation momentum





Idea: Since all BFKL paths count the same, just drop all paths going beyond $\rho_s(Y)$

Can just put in absorptive boundary to eliminate "bad" BFKL paths of evolution.



Dionysis T., A.M.

Solve BFKL with absorptive bndry.

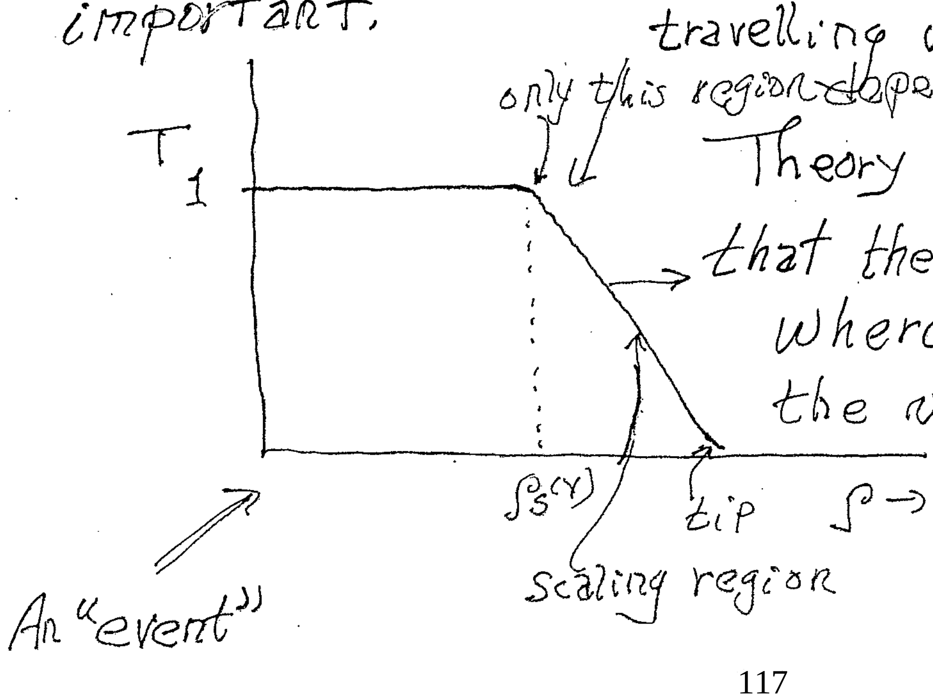
Gribov, Levin, Ryskin ('83); Golec-Biernat, Motyka, Stasto
Iancu, Itakura, McLerran

Find

$$\rho_s(Y) = \frac{2\alpha N_c}{\pi}\frac{\chi(\lambda_0)}{1-\lambda_0}Y - \frac{3}{2(1-\lambda_0)}\ln \alpha Y - \frac{2}{1-\lambda_0}\left[\ln\frac{1}{\alpha} - \frac{1}{2}\ln\ln\frac{1}{\alpha}\right] + C$$

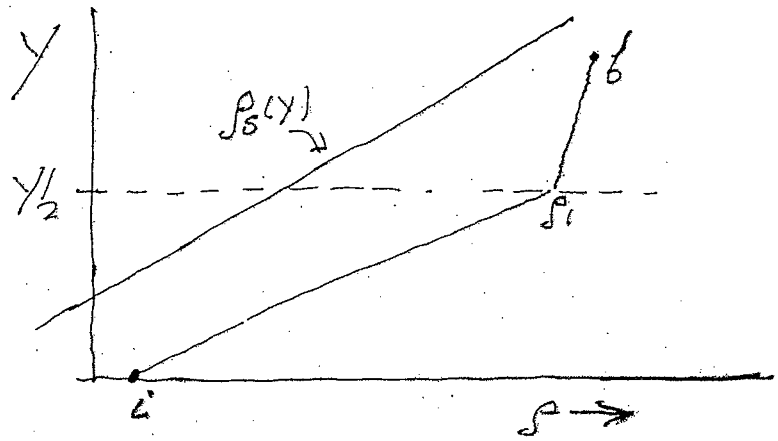
as result for Kovchegov equation.

Munier + Peschanski derive result in very general way. Kovchegov equation is of FKPP-type. Exact way in which unitarity imposed not important.



Theory of "pulled fronts" says that the tip of the front, where T is small, determines the velocity of the front, $\frac{d\rho_s}{dY}$.

(c) A problem with the Kovchegov equation A.Shoshi, A.M.

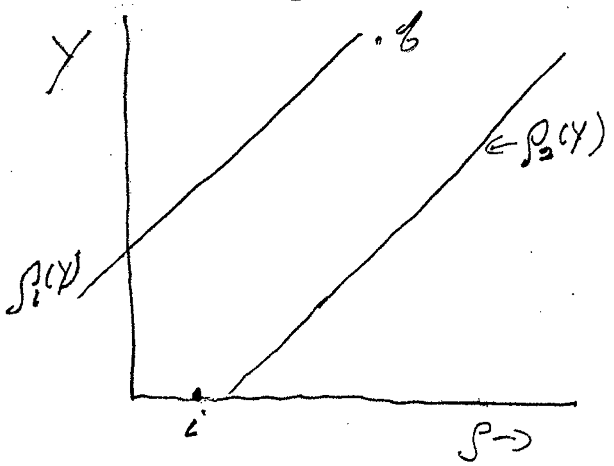


$$T(\rho_b,\rho_i,Y) \propto \frac{1}{\alpha^2}\int d\rho_1\, T(\rho_b,\rho_1,\tfrac{Y}{2})\, T(\rho_1,\rho_i,\tfrac{Y}{2})$$

Dominant contribution from

$$T(\rho_1,\rho_i,Y/2) \ll \alpha^2 \quad \underbrace{T(\rho_b,\rho_1,Y/2) \gg 1}_{\text{Violates unitarity!!}}$$

Introduce a second boundary to keep unitarity



Do BFKL evolution only in $\rho_1(Y) < \rho < \rho_2(Y)$

Find $$\rho_s(Y) = \frac{2\alpha N_c}{\pi}\frac{\chi(\lambda_0)}{1-\lambda_0}Y\left[1 - \frac{\pi^2 \chi''(\lambda_0)}{2(\Delta\rho)^2\chi(\lambda_0)}\right] + O\left(\frac{\alpha Y \ln\Delta\rho}{(\Delta\rho)^3}\right)$$

$$\Delta\rho = \rho_2 - \rho_1 = \frac{2}{1-\lambda_0}\ln 1/\alpha$$

Parametrically small correction, but ...

2. Relation to reaction-diffusion processes in statistical physics

E. Iancu, S. Munier, A.M.

The Brunet-Derrida model

Journal of Stat. Phys 103 (2001) 269.

N particles have positions $X_i(t)$. X_i is an integer and $t = 0, 1, 2, 3 \cdots$. Given $\{X_i(t)\}$, $X_k(t+1)$ is given by the rule: Choose $i, i' \leq N$ randomly, then

$$X_k(t+1) = \max[X_i(t), X_{i'}(t)] + \alpha_k(t+1)$$

where
$$\alpha_k(t+1) = \begin{cases} 0 & \text{with probability } 1-p \\ 1 & \text{with probability } p \end{cases}$$

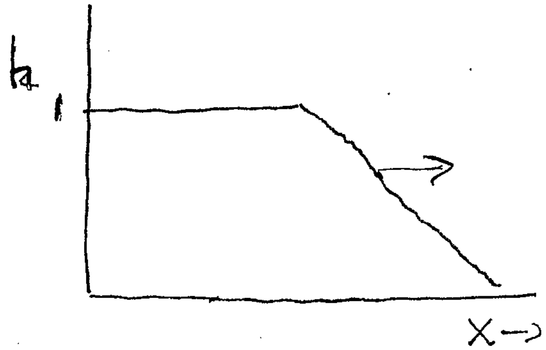
Define
$$h(x,t) = \frac{1}{N} \sum_{X_k(t) > X} 1$$

(We assume N large)

(a) The deterministic limit:

When N is very large we might expect that $h(x,t) \simeq \langle h(x,t) \rangle$. In this case

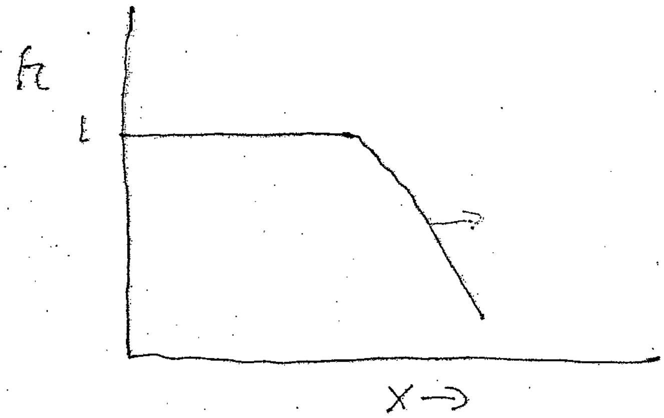
$$h(x,t+1) = 1 - p[1 - h(x-1,t)]^2 - (1-p)[1-h(x,t)]^2$$ $h = 0, 1$ are fixed points

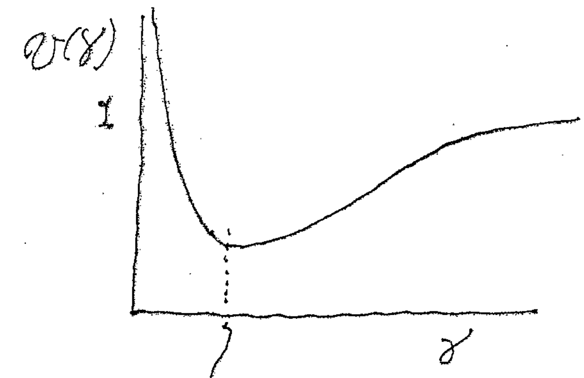


$$h(x,t) \sim e^{-\gamma(X - v(\gamma)t)}$$

$$v(\gamma) = \frac{1}{\gamma} \ln[2p e^{\gamma} + 2(1-p)]$$

analog of $\frac{\chi(\lambda)}{1-\lambda}$





velocity of Front $= v(\gamma^*)$

BD

$$e^{-\gamma^*(x - v(\gamma^*) t)}$$

BFKL

$$e^{2\alpha\chi(\lambda_0)Y - (1-\lambda_0)\rho} = e^{-(1-\lambda_0)\left[\rho - \frac{2\alpha\chi(\lambda_0)}{1-\lambda_0}Y\right]}$$

$$v(\gamma^*) \Longleftrightarrow \frac{2\alpha\chi(\lambda_0)}{1-\lambda_0}Y$$

The cutoff approximation:

When $n \sim 1/N$ the continuum approximation clearly makes no sense. BD set n to zero when the deterministic continuum equation reaches $1/N$. They Find

BD

$$v = v(\gamma^*)\left[1 - \frac{\pi^2 v''(\gamma^*)}{2\left[\frac{1}{\gamma^*}\ln N\right]^2 v(\gamma^*)}\right]$$

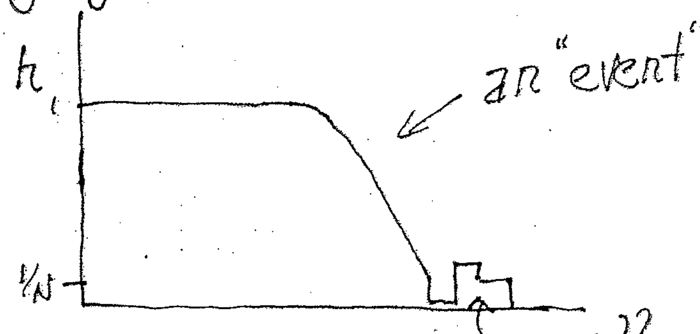
compare to AM, AS

$$\rho_s(Y) = \frac{2\alpha N_c}{\pi}\frac{\chi(\lambda_0)}{1-\lambda_0}Y\left[1 - \frac{\pi^2\chi''(\lambda_0)}{2(\Delta\rho)^2\chi(\lambda_0)}\right]$$

Identical results!

$N \sim 1/\alpha^2$

(b) Bringing back the stochasticity:



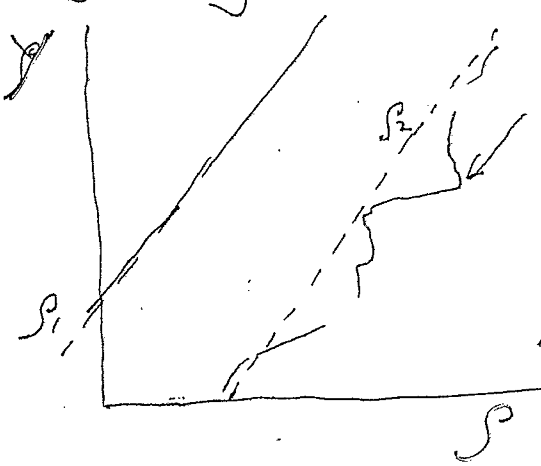
small, of size $1/N$, but big influence!

From numerical simulations by BD..., but converted to QCD variables

$$\langle \rho_s^2 \rangle - \langle \rho_s \rangle^2 = c \frac{\alpha Y}{(\Delta\rho)^3}$$

Should be a general result for all pulled fronts.

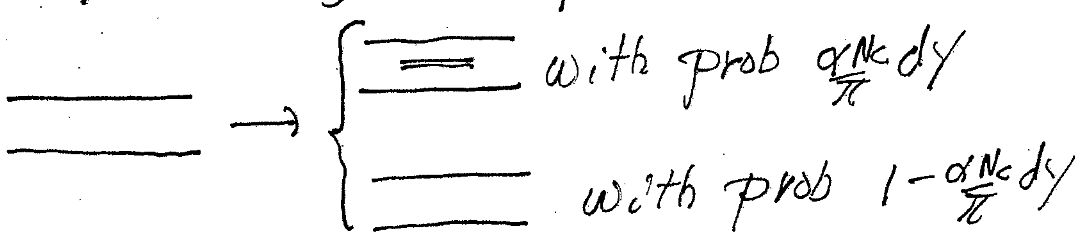
(c) Joining Forces E. Brunet, B. Derrida, S. Munier, A.M.



allowing fluctuations of front beyond ρ_2, by amount $\ln\Delta\rho$ create fluctuations of $\rho_s - \langle\rho_s\rangle$ of size $\sqrt{\frac{\alpha Y}{(\Delta\rho)^3}}$. Such fluctuations occur once in an interval $\Delta\alpha Y \simeq (\Delta\rho)^3$ and are of size $\rho = \rho_2 + \frac{3}{1-\gamma_0}\ln\Delta\rho$.

3. Back to QCD; Fluctuations and the event by event picture

Here the stochasticity is quantum mechanics. When dipole occupancy is low a dipole may or may not split in an interval dy

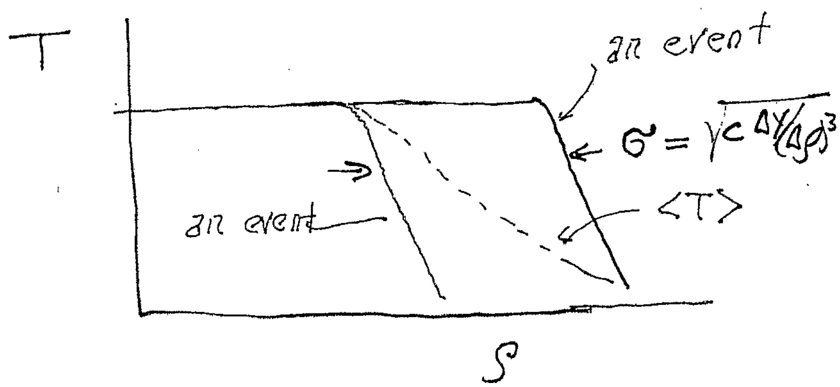


At high occupancy this stochasticity is not important, but can "pull" front from low occupancy.

$$\sigma^2 = \langle \rho_s^2 \rangle - \langle \rho_s \rangle^2 = c\,\frac{\alpha Y}{(\Delta\rho)^3}$$ naturally leads to

$$T(\rho, Y) = T\left(\frac{\rho - \langle \rho_s(Y) \rangle}{\sqrt{\alpha Y/(\Delta\rho)^3}} \right)$$

because the saturation momentum is not well defined up to size $\sqrt{\Delta Y/(\Delta\rho)^3}$.



Classical and Quantum Aspects of the Color Glass Condensate

March 7-11,2005

Registered Participants

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Classical and Quantum Aspects of the Color Glass Condensate

Physics Department, Bldg. 510
Large Seminar Room
March 7-11,2005

March 7, Monday

- 9:00-9:40 REGISTRATION
- 9:40-10:20 Larry McLerran, *Through the colored looking glass*
- 10:20-11:00 Kazunori Itakura, *Fluctuation-saturation duality andpomeron effective theoy*
- 11:00-11:30 COFFEE BREAK
- 11:30-12:10 Boris Kopeliovich, *Effects mocking CGC*
- 12:10-13:30 LUNCH
- 14:00-16:00 Discussion: Leader, Yuri Kovcliegov
- 18:30-20:30 WORKSHOP DINNER

March 8, Tuesday

- 9:00-9:40 Alfred Mueller, *Saturation, unitarity and fluctuations in high energy collisions*
- 9:40-10:20 Heribert Weigert, *Of colored glass and jets in a medium*
- 10:20-10:50 COFFEEBREAK
- 10:50-11:30 Yoshitaka Hatta, *Perturbative odderon in the color glass condensate*
- 11:30-12:10 Stephen Wong, *The extended evolution equation of the CGC in a dilute medium*
- 12:10-13:30 LUNCH
- 14:00-16:00 Discussion: Leader, Larry McLerran

March 9, Wednesday

- 9:00-9:40 Jochen Bartels, *Reggeonfield theoy in QCD*
- 9:40-10:20 Alex Kovner ,*B-JIMWLK^2: beyond and beside*
- 10:20-10:50 COFFEEBREAK
- 10:50-11:30 Francois Gelis, *Quarkproduction in pA collisions: rescattering effects and kt-factorization breaking*
- 11:30-12:10 Carlos Salgado, *The geometric scaling, the behavior of the saturation scale & the experimental data*
- 12:10-13:30 LUNCH
- 14:00-16:00 Discussion: Leader, Adrian Dumitru

March 10, Thursday

- 9:00-9:40 Ian Balitsky, *Scattering of shock waves in QCD*
- 9:40-10:20 Jianwei Qiu, *Transition from naiveparton model toparton saturation*
- 10:20-10:50 COFFEE BREAK
- 10:50-11:30 Elena Gonzalez Ferreiro, *Colour strings versus hardpomeron*
- 11:30-12:10 Michael Lublinsky, *New developments in the dipole model*
- 12:10-13:30 LUNCH
- 14:00-16:00 Discussion: Leader, Alfred Mueller

March 11, Friday

- 9:00-9:40 Yuri Kovchegov, *Can hydrodynamic description of heavy Ion collisions be derivedfvom Feynman diagrams?*
- 9:40-10:20 Ismail Zahed, *Strongly coupled QGP :reality from duality*
- 10:20-10:50 COFFEE BREAK
- 10:50-12:10 Adrian Dumitru, *Observational constraints on Qs from cosmic ray airshower data*
- 12:10-13:30 LUNCH
- 14:00-15:00 Discussion: Leader, Jochen Bartels
- 15:00-15:45 Workshop Summary

Additional RIKEN BNL Research Center Proceedings:

Additional RIKEN BNL Research Center Proceedings:

Volume 36 – RHIC Spin Collaboration Meeting VI – BNL-52642
Volume 35 – RIKEN Winter School – Quarks, Hadrons and Nuclei – QCD Hard Processes and the Nucleon Spin – BNL-52643
Volume 34 – High Energy QCD: Beyond the Pomeron – BNL-52641
Volume 33 – Spin Physics at RHIC in Year-1 and Beyond – BNL-52635
Volume 32 – RHIC Spin Physics V – BNL-52628
Volume 31 – RHIC Spin Physics III & IV Polarized Partons at High Q^2 Region – BNL-52617
Volume 30 – RBRC Scientific Review Committee Meeting – BNL-52603
Volume 29 – Future Transversity Measurements – BNL-52612
Volume 28 – Equilibrium & Non-Equilibrium Aspects of Hot, Dense QCD – BNL-52613
Volume 27 – Predictions and Uncertainties for RHIC Spin Physics & Event Generator for RHIC Spin Physics III – Towards Precision Spin Physics at RHIC – BNL-52596
Volume 26 – Circum-Pan-Pacific RIKEN Symposium on High Energy Spin Physics – BNL-52588
Volume 25 – RHIC Spin – BNL-52581
Volume 24 – Physics Society of Japan Biannual Meeting Symposium on QCD Physics at RIKEN BNL Research Center – BNL-52578
Volume 23 – Coulomb and Pion-Asymmetry Polarimetry and Hadronic Spin Dependence at RHIC Energies – BNL-52589
Volume 22 – OSCAR II: Predictions for RHIC – BNL-52591
Volume 21 – RBRC Scientific Review Committee Meeting – BNL-52568
Volume 20 – Gauge-Invariant Variables in Gauge Theories – BNL-52590
Volume 19 – Numerical Algorithms at Non-Zero Chemical Potential – BNL-52573
Volume 18 – Event Generator for RHIC Spin Physics – BNL-52571
Volume 17 – Hard Parton Physics in High-Energy Nuclear Collisions – BNL-52574
Volume 16 – RIKEN Winter School - Structure of Hadrons - Introduction to QCD Hard Processes – BNL-52569
Volume 15 – QCD Phase Transitions – BNL-52561
Volume 14 – Quantum Fields In and Out of Equilibrium – BNL-52560
Volume 13 – Physics of the 1 Teraflop RIKEN-BNL-Columbia QCD Project First Anniversary Celebration – BNL-66299
Volume 12 – Quarkonium Production in Relativistic Nuclear Collisions – BNL-52559
Volume 11 – Event Generator for RHIC Spin Physics – BNL-66116
Volume 10 – Physics of Polarimetry at RHIC – BNL-65926
Volume 9 – High Density Matter in AGS, SPS and RHIC Collisions – BNL-65762
Volume 8 – Fermion Frontiers in Vector Lattice Gauge Theories – BNL-65634
Volume 7 – RHIC Spin Physics – BNL-65615
Volume 6 – Quarks and Gluons in the Nucleon – BNL-65234
Volume 5 – Color Superconductivity, Instantons and Parity won?)-Conservation at High Baryon Density – BNL-65105

Additional RIKEN BNL Research Center Proceedings:

Volume 4 – Inauguration Ceremony, September 22 and Non -Equilibrium Many Body Dynamics – BNL-64912

Volume 3 – Hadron Spin-Flip at RHIC Energies – BNL-64724

Volume 2 – Perturbative QCD as a Probe of Hadron Structure – BNL-64723

Volume 1 – Open Standards for Cascade Models for RHIC – BNL-64722

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RIKEN BNL RESEARCH CENTER

Classical and Quantum Aspects of the Color Glass Condensate

March 7-11,2005

Li Keran

Nuclei as heavy as bulls
Through collision
Generate new states of matter.
T.D. Lee

Speakers:

Ian Balitsky	Jochen Bartels	Adrian Dumitru	Elena Gonzalez Ferreiro
Francois Gelis	Yoshitaka Hatta	Kazunori Itakura	Boris Kopeliovich
Yuri Kovchegov	Alex Kovner.	Michael Lublinsky	Larry McLerran
Alfred Mueller	Jianwei Qiu	Carlos Salgado	Heribert Weigert
Stephen Wong	Ismail Zahed		

Organizers: D. Kharzeev, A. Stasto, K. Tuchin, W. Vogelsang